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On One-loop Neutron Electric Dipole Moment from Supersymmetry without R-parity

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Abstract

We present a detailed analysis together with exact numerical calculations on one-loop contributions to neutron electric dipole moment from supersymmetry without R-parity, focusing on the gluino, chargino, and neutralino contributions. Apart from the neglected family mixing among quarks, complete formulae are given for the various contributions, through the quark dipole operators, to which the present study is restricted. We discuss the structure and main features of the R-parity violating contributions and the interplay between the R-parity conserving and violating parameters. In particular, the parameter combination $\mu_i^* \lambda'_{i11}$, under the optimal parametrization adopted, is shown to be solely responsible for the R-parity violating contributions in the supersymmetric loop diagrams. While $\mu_i^* \lambda_{i11}$ could bear a complex phase, the latter is not necessary to have a R-parity violating contribution.

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I. INTRODUCTION

Neutron and electron electric dipole moments (EDM's) are important topics for new CP violating physics. They are known to be extremely small in the Standard Model (SM); in fact, way below the present experimental limit. With supersymmetry (SUSY) comes many plausible extra EDM contributions. That has led to the so called SUSY-CP problem [1] for the minimal supersymmetric standard model (MSSM). If one simply takes the minimal supersymmetric spectrum of the SM and imposes nothing more than the gauge symmetries while still admits soft SUSY breaking, the generic supersymmetric standard model would be resulted. When the large number of baryon and/or lepton number violating terms in such a generic supersymmetric SM is removed by hand, through imposing an *ad hoc* discrete symmetry called R-parity, one obtains the MSSM Lagrangian. In the case of R-parity violation, two recent papers focus on the contributions from the extra trilinear terms in the superpotential and conclude that there is no new EDM contribution at the 1-loop level [2]. Perhaps it has not been emphasized enough in the two papers that they are *not* studying the complete theory of SUSY without R-parity, which is nothing other than the generic supersymmetric SM; in particular, they have neglected admissible RPV parameters other than the trilinear ones in the superpotential. It is interesting to see that in the generic case there is in fact contribution at 1-loop level, as pointed out in Refs. [3,4]. In particular, Ref. [3] gives a clear illustration of the much overlooked existence of a R-parity violating (RPV) contribution to LR squark mixings and the resulted contribution to neutron EDM through the simple 1-loop gluino diagram. We would like to emphasize again that the new contribution involves both bilinear and trilinear (RPV) couplings in the superpotential. Since RPV scenarios studied in the literature typically admit only one of the two types of couplings, the contribution has not been previously identified. A simple estimate on the bound obtained on the RPV parameters (the $\mu_i^* \lambda'_{i11}$ combination) given in Ref. [3] has already illustrated that the bound from the neutron EDM as one of the most important, being competitive even when compared with sub-eV neutrino mass bounds and substantially more stringent than most collider bounds. The present article aims at giving a detailed analysis and numerical study of the RPV extension of SUSY contributions to neutron EDM. Similar new RPV contributions to electron EDM have been noted in Ref. [3]. In fact, the complete result for RPV contributions to the masses of the sleptons and other (color-singlet) scalars has been given in Ref. [5], which focuses mainly on their implications to neutrino masses.

The complete theory of supersymmetry (SUSY) without R-parity admits all kind of RPV terms without bias. It is obviously better motivated than *ad hoc* versions of RPV theories. The large number of new parameters involved, however, makes the theory difficult to analyze. The question of the specification of flavor bases to define the parameters in the

Lagrangian of the theory unambiguously becomes more important. In fact, thinking about the theory as the generic supersymmetric SM instead of as “MSSM + RPV terms” helps to clarify many of the issues involved [6]. From such a perspective, it has been illustrated [7] that an optimal parametrization, called the single-VEV parametrization (SVP), provides a very nice formulation which helps to simplify much of the analysis. In particular, the SVP gives the complete results for the tree-level mass matrices of all state, fermions as well as scalars, in the simplest form [5]. The formulation has been used to study leptonic phenomenology [7] and various aspect of neutrino masses [5,8–11]. The present EDM study (also Refs. [3,4]), and the parallel works on $\mu \rightarrow e \gamma$ [12] (see also Ref. [13]), electron EDM and $b \rightarrow s \gamma$ [14] will further illustrate the advantage of adopting the SVP.

We focus here only on such contributions to the neutron EDM, based on the valence quark model [15]. Hence, we study only the 1-loop quark EDM’s. We will give complete 1-loop formulae for EDM’s of the up- and down-sector quarks, of which the u and d results are used to calculate the neutron EDM through the

$$d_n = \frac{1}{3} (4 d_d - d_u) \eta \quad (1)$$

formula, where $\eta \simeq 1.53$ is a QCD correction factor from renormalization group evolution [16,17]. This is to be matched with the experimental bound [18] :

$$d_n < 6.3 \cdot 10^{-26} e \cdot \text{cm} .$$

In the MSSM case, one has the SUSY loop contributions and the charged Higgs contributions. The latter is very negligible. We focus here in this article on the analog of the SUSY loop contributions. The latter includes the gluino loop, the chargino-like loop, and the neutralino-like loop. By the last two, we mean generalization of the chargino and neutralino loops under the generic picture. The (RPV) mixings of the leptons with the gauginos and higgsinos gives five (color-singlet) charged leptons and seven neutral fermions, including the charginos and neutralinos as well as e , μ , τ , and three physical neutrinos. They come from the same set of electroweak states and should not be separated from one another in the analysis. It is no surprise that the physical chargino and neutralino states dominate the EDM contributions. We use explicit exact mass eigenstate expressions in our analysis to illustrate that as well as other interesting features, starting from the generic electroweak states couplings under the SVP. An exact numerical calculation is performed. We would like to mention that the generalization of the charged Higgs loop contribution involves other different RPV parameters. Moreover, there are many new and potentially important contributions including a t quark loop, as also pointed out in Ref. [4]. We will give also the formulae of such quark-scalar loop contributions, though a detailed study is postponed to a future publication. Note that Ref. [4], which first appeared around the same time as Ref. [3], is the only other study on the same topic available. In our opinion, our study here is

more systematic and complete. Ref. [4] does not include, for example, the RPV LR scalar mixing and the resulted gaugino loop contribution to EDM. Moreover, to the best of our knowledge, the present study includes the first exact numerical calculation performed. Ref. [4] also quotes a $\cos\beta$ dependence of the major chargino-like contribution, hence weakening of bound in large $\tan\beta$ regime — a result we disagree. Our careful numerical study illustrates many more interesting features, as the discussion below will speak for itself.

This paper is organized as follows: We first summarize the formulation and notation used in Sec. II, where we also elaborate in some detail the electroweak fermion field couplings needed to study the quark EDM's. Next, Sec. III contains results presented recently in Ref. [5] on all the scalar masses, listed here so that the present paper will be self-contained. Of most importance here is the RPV contributions to LR mixings, which play a central part in the EDM contributions. The quark EDM's are analyzed in Sec. IV. Some results from our numerical calculations are presented in Sec. V, after which we conclude in Sec. VI. An appendix gives some background formulae on the color-singlet fermion masses. Note that our formulae and calculations here naturally include the R-parity conserving MSSM part, though we will concentrate on the role of the RPV parameters and their unique contributions. Corresponding studies of the MSSM case can be found, for example, in Refs. [16,19–21], to which we refer the readers for comparison.

II. FORMULATION AND NOTATION

We summarize our formulation and notation below. The most general renormalizable superpotential for the generic supersymmetric SM can be written then as

$$W = \varepsilon_{ab} \left[\mu_\alpha \hat{H}_u^a \hat{L}_\alpha^b + y_{ik}^u \hat{Q}_i^a \hat{H}_u^b \hat{U}_k^c + \lambda'_{\alpha jk} \hat{L}_\alpha^a \hat{Q}_j^b \hat{D}_k^c + \frac{1}{2} \lambda_{\alpha\beta k} \hat{L}_\alpha^a \hat{L}_\beta^b \hat{E}_k^c \right] + \frac{1}{2} \lambda''_{ijk} \hat{U}_i^c \hat{D}_j^c \hat{D}_k^c, \quad (2)$$

where (a, b) are $SU(2)$ indices, (i, j, k) are the usual family (flavor) indices, and (α, β) are extended flavor index going from 0 to 3. At the limit where $\lambda_{ijk}, \lambda'_{ijk}, \lambda''_{ijk}$ and μ_i all vanish, one recovers the expression for the R-parity preserving case (*i.e.* MSSM), with \hat{L}_0 identified as \hat{H}_d . Without R-parity imposed, the latter is not *a priori* distinguishable from the \hat{L}_i 's. Note that λ is antisymmetric in the first two indices, as required by the $SU(2)$ product rules, as shown explicitly here with $\varepsilon_{12} = -\varepsilon_{21} = 1$. Similarly, λ'' is antisymmetric in the last two indices from $SU(3)_C$.

R-parity is exactly an *ad hoc* symmetry put in to make \hat{L}_0 , stand out from the other \hat{L}_i 's as the candidate for \hat{H}_d . It is defined in terms of baryon number, lepton number, and spin as, explicitly, $\mathcal{R} = (-1)^{3B+L+2S}$. The consequence is that the accidental symmetries of baryon number and lepton number in the SM are preserved, at the expense of making particles and superparticles having a categorically different quantum number, R-parity. The

latter is actually not the most effective discrete symmetry to control superparticle mediated proton decay [22], but is most restrictive for in terms of what is admitted in the Lagrangian, or the superpotential alone.

A naive look at the scenario suggests that the large number of new couplings makes the task formidable. However, it becomes quite manageable with an optimal choice of flavor bases, the SVP [7]. In fact, doing phenomenological studies without specifying a choice of flavor bases is ambiguous. It is like doing SM quark physics with 18 complex Yukawa couplings instead of the 10 real physical parameter. For SUSY without R-parity, the choice of an optimal parametrization mainly concerns the 4 \hat{L}_α flavors. Under the SVP¹, flavor bases are chosen such that : 1/ among the \hat{L}_α 's, only \hat{L}_0 , bears a VEV (*i.e.* $\langle \hat{L}_i \rangle \equiv 0$); 2/ $y_{jk}^e (\equiv \lambda_{0jk} = -\lambda_{j0k}) = \frac{\sqrt{2}}{v_0} \text{diag}\{m_1, m_2, m_3\}$; 3/ $y_{jk}^d (\equiv \lambda_{0jk}) = \frac{\sqrt{2}}{v_0} \text{diag}\{m_d, m_s, m_b\}$; 4/ $y_{ik}^u = \frac{v_u}{\sqrt{2}} V_{\text{CKM}}^T \text{diag}\{m_u, m_c, m_t\}$, where $v_0 \equiv \sqrt{2} \langle \hat{L}_0 \rangle$ and $v_u \equiv \sqrt{2} \langle \hat{H}_u \rangle$. The big advantage of here is that the (tree-level) mass matrices for all the fermions do not involve any of the trilinear RPV couplings, though the approach makes no assumption on any RPV coupling including even those from soft SUSY breaking; and all the parameters used are uniquely defined, with the exception of some removable phases.

We are interested in scalar-fermion-fermion couplings similar to those of the charginos and neutralinos in MSSM. The gaugino couplings are, of course, standard. Coming from the gauge interaction parts, they have nothing to do with R-parity. The “higgsino-like” part is, however, different. Without R-parity and in a generic flavor basis of the 4 \hat{L}_α 's, the \hat{H}_d of MSSM is hidden among the latter. The SVP however identify \hat{L}_0 as the one having “Higgs” properties of \hat{H}_d , though it still maintains (RPV) couplings similar to those of the \hat{L}_i 's. We write the components of a \hat{L}_α fermion doublet as l_α^0 and l_α^- , and their scalar partners as \tilde{l}_α^0 and \tilde{l}_α^- . Apart from being better motivated theoretically, the common notation helps to trace the flavor structure. However, we will also use notations of the form h_d^\star and \tilde{h}_d^\star , as alternative notations for \tilde{l}_0^\star and l_0^\star , in some places below. This is unambiguous under our formulation. We will also referred to h_d^\star and \tilde{h}_d^\star as Higgs and higgsino, respectively; while they are generally also included in the terms slepton and lepton.

Note that in the left-handed lepton and slepton field notations introduced above, we have dropped the commonly used L -subscript, for simplicity. For the components of the three right-handed leptonic superfields, we use l_i^+ and \tilde{l}_i^+ , with again the R -subscript dropped. The notation for the quark and squark fields will be standard, with the L - and R -subscripts. A normal quark state, such as d_{Lk} , denotes a mass eigenstate, while a squark state the

¹Note that our notations here are a bit different from those in the reference; we follow mostly the notations in Refs. [3] and [5] while improving and elaborating further whenever appropriate. We will clarify all notations used as our discussion goes along [6].

supersymmetric partner of one. A quark or squark state with a \prime denotes one with the quark state being the $SU(2)$ partner of a mass eigenstate. For instance, \tilde{u}'_{L_3} is the up-type squark state from \hat{Q}_3 which contains the exact left-handed b quark according to our parametrization of the Lagrangian.

The up-sector Higgs is unaffected by R-parity violation. The scalar and fermion states of the doublet are denoted by h_u^+ and h_u^0 , and \tilde{h}_u^+ and \tilde{h}_u^0 , respectively.

The identity of the charginos and neutralinos is unambiguous in MSSM. Without R-parity, they mix with the charged leptons and neutrinos. In fact, the true charged leptons and neutrinos are the light mass eigenstates of the full 5×5 and 7×7 mass matrices respectively. The mass eigenstates deviate from the l_i^- 's and l_i^0 's. Though the latter deviations are practically negligible in the limit of small μ_i 's, it still helps to distinguish the electroweak states from the mass eigenstates. Moreover, large R-parity violation, especially in terms of a large μ_3 , is not definitely ruled out [7]. Here in this paper, we make the explicit distinction and reserve the terms chargino and neutralino for the heavy states beyond the physical charged leptons e , μ and τ and neutrinos. The m_i 's introduced above are the Yukawa contributions to the physical charged lepton masses, hence approximately equal to the latter. The readers are referred to appendix A for more details about the fermion mass terms.

We are now ready to spell out the couplings concerning the (color- singlet) charged and neutral fermion from the superpotential. We have

$$\begin{aligned} \mathcal{L}_\chi = & y_{u_i} V_{\text{CKM}}^{ij} \tilde{h}_u^+ \left[\tilde{u}_{R_i}^c d_{L_j} + u_{R_i}^c \tilde{d}_{L_j} \right] + y_{d_i} l_0^- \left[\tilde{d}_{R_i}^c u'_{L_i} + d_{R_i}^c \tilde{u}'_{L_i} \right] + \lambda_{ijk} l_i^- \left[\tilde{d}_{R_k}^c u'_{L_j} + d_{R_k}^c \tilde{u}'_{L_j} \right] \\ & - y_{u_i} \tilde{h}_u^0 \left[\tilde{u}_{R_i}^c u_{L_i} + u_{R_i}^c \tilde{u}_{L_i} \right] - y_{d_i} l_0^0 \left[\tilde{d}_{R_i}^c d_{L_i} + d_{R_i}^c \tilde{d}_{L_i} \right] - \lambda_{ijk} l_i^0 \left[\tilde{d}_{R_k}^c d_{L_j} + d_{R_k}^c \tilde{d}_{L_j} \right] \\ & + y_{e_i} \left[l_0^- l_i^+ \tilde{l}_i^0 - l_i^- l_i^+ \tilde{l}_0^0 \right] + y_{e_i} \left[l_0^- l_i^0 \tilde{l}_i^+ - l_i^- l_0^0 \tilde{l}_i^+ \right] + y_{e_i} \left[l_i^0 l_i^+ \tilde{l}_0^- - l_0^0 l_i^+ \tilde{l}_i^- + \right. \\ & \left. + \lambda_{ijk} l_i^- l_k^+ \tilde{l}_j^0 + \lambda_{ijk} l_i^- l_j^0 \tilde{l}_k^+ - \lambda_{ijk} l_i^0 l_k^+ \tilde{l}_j^- - \lambda_{ijk} l_i^0 l_k^+ \tilde{l}_j^- \right] + \text{h.c.} , \end{aligned} \quad (3)$$

where

$$y_{u_i} = \frac{g_2 m_{u_i}}{\sqrt{2} M_W \sin\beta} , \quad y_{d_i} = \frac{g_2 m_{d_i}}{\sqrt{2} M_W \cos\beta} , \quad y_{e_i} = \frac{g_2 m_i}{\sqrt{2} M_W \cos\beta} , \quad (4)$$

are the (diagonal) quark and charged lepton Yukawa couplings; and $\tan\beta = \frac{v_u}{v_0}$ [23]. Recall that λ_{0jk} corresponds to the down-Yukawa coupling matrix, and λ_{0jk} corresponds to the charged lepton Yukawa coupling matrix, both of which are diagonal under the SVP; in addition, we have $u'_{L_i} = V_{\text{CKM}}^{\dagger ij} u_{L_j}$ being $SU(2)$ partner of the mass eigenstate d_{L_i} , and \tilde{u}'_{L_i} its scalar partner. We also use below \tilde{d}'_{L_i} , which is, explicitly, $V_{\text{CKM}}^{ij} \tilde{d}_{L_j}$. There are some more scalar-fermion-fermion terms besides those given in \mathcal{L}_χ . These extra terms are slepton-quark-quark terms. We will see that the latter, is actually also involved in 1-loop EDM diagrams, though not the major focus of this paper. With the above explicit listed terms, however, it is straight forward to see what the extra terms are like.

In both of the above expressions for \mathcal{L}_χ , there is clear distinction between the MSSM terms and the RPV terms. The nice feature is a consequence of the SVP. The simple structure of the trilinear coupling contributions to the d -quark and charged lepton masses, which is equivalent to that of the R-parity conserving limit, is what make the analysis simple and easy to handle. We want to emphasize that the above expressions are exact tree-level results without hidden assumptions behind its validity. They are good even when there is large R-parity violation.

III. SQUARK AND SLEPTON MASSES

The soft SUSY breaking part of the Lagrangian, in terms of scalar parts of the superfield multiplets, can be written as follows :

$$\begin{aligned}
V_{\text{soft}} = & \epsilon_{ab} B_\alpha H_u^a \tilde{L}_\alpha^b + \epsilon_{ab} \left[A_{ij}^U \tilde{Q}_i^a H_u^b \tilde{U}_j^c + A_{ij}^D H_d^a \tilde{Q}_i^b \tilde{D}_j^c + A_{ij}^E H_d^a \tilde{L}_i^b \tilde{E}_j^c \right] + \text{h.c.} \\
& + \epsilon_{ab} \left[A_{ijk}' \tilde{L}_i^a \tilde{Q}_j^b \tilde{D}_k^c + \frac{1}{2} A_{ijk}'' \tilde{L}_i^a \tilde{L}_j^b \tilde{E}_k^c \right] + \frac{1}{2} A_{ijk}''' \tilde{U}_i^c \tilde{D}_j^c \tilde{D}_k^c + \text{h.c.} \\
& + \tilde{Q}^\dagger \tilde{m}_Q^2 \tilde{Q} + \tilde{U}^\dagger \tilde{m}_U^2 \tilde{U} + \tilde{D}^\dagger \tilde{m}_D^2 \tilde{D} + \tilde{L}^\dagger \tilde{m}_L^2 \tilde{L} + \tilde{E}^\dagger \tilde{m}_E^2 \tilde{E} + \tilde{m}_{H_u}^2 |H_u|^2 \\
& + \frac{M_1}{2} \tilde{B} \tilde{B} + \frac{M_2}{2} \tilde{W} \tilde{W} + \frac{M_3}{2} \tilde{g} \tilde{g} + \text{h.c.} ,
\end{aligned} \tag{5}$$

where we have separated the R-parity conserving ones from the RPV ones ($H_d \equiv \tilde{L}_0$) for the A -terms. Note that $\tilde{L}^\dagger \tilde{m}_L^2 \tilde{L}$, unlike the other soft mass terms, is given by a 4×4 matrix. Explicitly, $\tilde{m}_{L_{00}}^2$ is $\tilde{m}_{H_d}^2$ of the MSSM case while $\tilde{m}_{L_{0k}}^2$'s give RPV mass mixings. The other notations are obvious.

The SVP also simplifies much the otherwise extremely complicated expressions for the mass-squared matrices of the scalar sectors. Firstly, we will look at the squarks sectors. The masses of up-squarks obviously have no RPV contribution. The down-squark sector, however, has interesting result. We have the mass-squared matrix as follows :

$$\mathcal{M}_D^2 = \begin{pmatrix} \mathcal{M}_{LL}^2 & \mathcal{M}_{RL}^{2\dagger} \\ \mathcal{M}_{RL}^2 & \mathcal{M}_{RR}^2 \end{pmatrix} , \tag{6}$$

where

$$\begin{aligned}
\mathcal{M}_{LL}^2 &= \tilde{m}_Q^2 + m_D^\dagger m_D + M_Z^2 \cos 2\beta \left[-\frac{1}{2} + \frac{1}{3} \sin^2 \theta_W \right] , \\
\mathcal{M}_{RR}^2 &= \tilde{m}_D^2 + m_D m_D^\dagger + M_Z^2 \cos 2\beta \left[-\frac{1}{3} \sin^2 \theta_W \right] ,
\end{aligned} \tag{7}$$

and

$$\begin{aligned}
(\mathcal{M}_{RL}^2)^T &= A^D \frac{v_0}{\sqrt{2}} - (\mu_\alpha^* \chi_{\alpha j k}) \frac{v_u}{\sqrt{2}} \\
&= [A_d - \mu_0^* \tan \beta] m_D + \frac{\sqrt{2} M_W \cos \beta}{g_2} \delta A^D - \frac{\sqrt{2} M_W \sin \beta}{g_2} (\mu_i^* \chi_{i j k}) .
\end{aligned} \tag{8}$$

Here, m_D is the down-quark mass matrix, which is diagonal under the parametrization adopted; A_d is a constant (mass) parameter representing the “proportional” part of the A -term and the matrix δA^p is the “proportionality” violating part; $(\mu_i^* \lambda_{ijk})$, and similarly $(\mu_\alpha^* \lambda_{\alpha jk})$, denotes the 3×3 matrix $()_{jk}$ with elements listed ². The $(\mu_\alpha^* \lambda_{\alpha jk})$ term is the full F -term contribution, while the $(\mu_i^* \lambda_{ijk})$ part separated out in the last expression gives the RPV contributions.

Next we go on to the slepton sector. From Eq.(5) above, we can see that the charged Higgses should be considered on the same footing together with the sleptons. We have hence an 8×8 mass-squared matrix of the following $1 + 4 + 3$ form :

$$\mathcal{M}_E^2 = \begin{pmatrix} \widetilde{\mathcal{M}}_{Hu}^2 & \widetilde{\mathcal{M}}_{LH}^{2\dagger} & \widetilde{\mathcal{M}}_{RH}^{2\dagger} \\ \widetilde{\mathcal{M}}_{LH}^2 & \widetilde{\mathcal{M}}_{LL}^2 & \widetilde{\mathcal{M}}_{RL}^{2\dagger} \\ \widetilde{\mathcal{M}}_{RH}^2 & \widetilde{\mathcal{M}}_{RL}^2 & \widetilde{\mathcal{M}}_{RR}^2 \end{pmatrix} ; \quad (9)$$

where

$$\begin{aligned} \widetilde{\mathcal{M}}_{Hu}^2 &= \tilde{m}_{Hu}^2 + \mu_\alpha^* \mu_\alpha + M_Z^2 \cos 2\beta \left[\frac{1}{2} - \sin^2 \theta_W \right] \\ &\quad + M_Z^2 \sin^2 \beta [1 - \sin^2 \theta_W] , \\ \widetilde{\mathcal{M}}_{LL}^2 &= \tilde{m}_L^2 + m_L^\dagger m_L + (\mu_\alpha^* \mu_\beta) + M_Z^2 \cos 2\beta \left[-\frac{1}{2} + \sin^2 \theta_W \right] , \\ &\quad + \begin{pmatrix} M_Z^2 \cos^2 \beta [1 - \sin^2 \theta_W] & 0_{1 \times 3} \\ 0_{3 \times 1} & 0_{3 \times 3} \end{pmatrix} , \\ \widetilde{\mathcal{M}}_{RR}^2 &= \tilde{m}_E^2 + m_E m_E^\dagger + M_Z^2 \cos 2\beta [-\sin^2 \theta_W] ; \end{aligned} \quad (10)$$

and

$$\begin{aligned} \widetilde{\mathcal{M}}_{LH}^2 &= (B_\alpha^*) + \begin{pmatrix} \frac{1}{2} M_Z^2 \sin 2\beta [1 - \sin^2 \theta_W] \\ 0_{3 \times 1} \end{pmatrix} , \\ \widetilde{\mathcal{M}}_{RH}^2 &= -(\mu_i^* \lambda_{i0k}) \frac{v_0}{\sqrt{2}} = (\mu_k^* m_k) \quad (\text{no sum over } k) , \\ (\widetilde{\mathcal{M}}_{RL}^2)^T &= \begin{pmatrix} 0 \\ A^E \end{pmatrix} \frac{v_0}{\sqrt{2}} - (\mu_\alpha^* \lambda_{\alpha \beta k}) \frac{v_u}{\sqrt{2}} \\ &= [A_e - \mu_0^* \tan \beta] \begin{pmatrix} 0 \\ m_E \end{pmatrix} + \frac{\sqrt{2} M_W \cos \beta}{g_2} \begin{pmatrix} 0 \\ \delta A^E \end{pmatrix} - \begin{pmatrix} -\mu_k^* m_k \tan \beta \\ \frac{\sqrt{2} M_W \sin \beta}{g_2} (\mu_i^* \lambda_{ijk}) \end{pmatrix} . \end{aligned} \quad (11)$$

Notations and results here are similar to the squark case above, with some difference. We have A_e and δA^E , or the extended matrices $\begin{pmatrix} 0 \\ * \end{pmatrix}$ incorporating them, denote the splitting of the A -term, with proportionality defined with respect to m_E ; $m_L = \text{diag}\{0, m_E\} =$

²Note that we use this kind of bracket notations for matrices extensively here. In this case, the repeated index i is to be summed over as usual, and hence dummy.

$\text{diag}\{0, m_1, m_2, m_3\}$. Recall that the m_i 's are approximately the charged lepton masses. A 4×3 matrix $(\mu_i^* \lambda_{i\beta k})$ gives the RPV contributions to $(\widetilde{\mathcal{M}}_{RL}^2)^T$. In the above expression, we separate explicitly the first row of the former, which corresponds to mass-squared terms of the type $\tilde{l}^+ h_d^-$ type ($h_d^- \equiv \tilde{l}_0^-$). The nonzero $\widetilde{\mathcal{M}}_{RH}^2$ and the B_i^* 's in $\widetilde{\mathcal{M}}_{LH}^2$ are also RPV contributions. The former is a $\tilde{l}^+(h_u^+)^{\dagger}$ type, while the latter a $\tilde{l}^- h_u^+$ term. Note that the parts with the $[1 - \sin^2 \theta_w]$ factor are singled out as they are extra contributions to the masses of the charged-Higgses (*i.e.* $l_0^- \equiv h_d^-$ and h_u^+). The latter is the result of the quartic terms in the scalar potential and the fact that the Higgs doublets bear VEV's.

The neutral scalar mass terms, in terms of the $(1+4)$ complex scalar fields, ϕ_n 's, can be written in two parts — a simple $(\mathcal{M}_{\phi\phi\ddagger}^2)_{mn} \phi_m^{\dagger} \phi_n$ part, and a Majorana-like part in the form $\frac{1}{2} (\mathcal{M}_{\phi\phi}^2)_{mn} \phi_m \phi_n + \text{h.c.}$. As the neutral scalars are originated from chiral doublet superfields, the existence of the Majorana-like part is a direct consequence of the electroweak symmetry breaking VEV's, hence restricted to the scalars playing the Higgs role only. They come from the quartic terms of the Higgs fields in the scalar potential. We have explicitly

$$\mathcal{M}_{\phi\phi}^2 = \frac{1}{2} M_Z^2 \begin{pmatrix} \sin^2 \beta & -\cos \beta \sin \beta & 0_{1 \times 3} \\ -\cos \beta \sin \beta & \cos^2 \beta & 0_{1 \times 3} \\ 0_{3 \times 1} & 0_{3 \times 1} & 0_{3 \times 3} \end{pmatrix} ; \quad (12)$$

and

$$\mathcal{M}_{\phi\phi\ddagger}^2 = \mathcal{M}_{\phi\phi}^2 + \begin{pmatrix} \tilde{m}_{H_u}^2 + \mu_{\alpha}^* \mu_{\alpha} + M_Z^2 \cos 2\beta \left[-\frac{1}{2}\right] & -(B_{\alpha}) \\ -(B_{\alpha}^*) & \tilde{m}_L^2 + (\mu_{\alpha}^* \mu_{\beta}) + M_Z^2 \cos 2\beta \left[\frac{1}{2}\right] \end{pmatrix} , \quad (13)$$

Note that $\mathcal{M}_{\phi\phi}^2$ here is real (see the appendix), while $\mathcal{M}_{\phi\phi\ddagger}^2$ does have complex entries. The full 10×10 (real and symmetric) mass-squared matrix for the real scalars is then given by

$$\mathcal{M}_S^2 = \begin{pmatrix} \mathcal{M}_{SS}^2 & \mathcal{M}_{SP}^2 \\ (\mathcal{M}_{SP}^2)^T & \mathcal{M}_{PP}^2 \end{pmatrix} , \quad (14)$$

where the scalar, pseudo-scalar, and mixing parts are

$$\begin{aligned} \mathcal{M}_{SS}^2 &= \text{Re}(\mathcal{M}_{\phi\phi\ddagger}^2) + \mathcal{M}_{\phi\phi}^2 , \\ \mathcal{M}_{PP}^2 &= \text{Re}(\mathcal{M}_{\phi\phi\ddagger}^2) - \mathcal{M}_{\phi\phi}^2 , \\ \mathcal{M}_{SP}^2 &= -2 \text{Im}(\mathcal{M}_{\phi\phi\ddagger}^2) , \end{aligned} \quad (15)$$

respectively. If $\text{Im}(\mathcal{M}_{\phi\phi\ddagger}^2)$ vanishes, the scalars and pseudo-scalars decouple from one another and the unphysical Goldstone mode would be found among the latter. Finally, we note that the B_{α} entries may also be considered as a kind of LR mixings.

We would like to emphasize that the above scalar mass results are complete — all RPV contributions, SUSY breaking or otherwise, are included. The simplicity of the result is

a consequence of the SVP. Explicitly, there are no RPV A -term contributions due to the vanishing of VEV's $v_i \equiv \sqrt{2}\langle\hat{L}_i\rangle$. The Higgs-slepton results given as in Eqs.(9) and (14) contain redundancy of parameters and hide the unphysical Goldstone state. However, the results as they are given here are good enough for some purposes including the present EDM discussion.

Note that in the following discussion, we will not consider flavor changing scalar mass terms from soft SUSY breaking, which could be suppressed through a flavor-blind SUSY breaking mediating mechanism such as gauge-mediation. A major concentration of our study, however, will be on the effects of the flavor changing scalar mass terms from RPV superpotential parameters, which give interesting new results.

IV. CONTRIBUTIONS TO NEUTRON EDM AT 1-LOOP

In this section, we will discuss the RPV 1-loop contributions to neutron EDM systematically, drawing comparison with the R-parity conserving (MSSM) case wherever it would be useful. We will give complete 1-loop formulae for quark EDM's for our generic supersymmetric SM. We will not, however, go into the chromoelectric dipole operator or Weinberg gluonic operator contributions. Following the common practice, family mixing will largely be neglected, though we will comment on some particularly interesting aspects of the issue. We will also compare the results with Feynman diagrams given more or less in electroweak basis. Naively, such diagrams, with minimal number of mass insertions admitted, should represent a first order result, at least where mass mixings are small. Note that the latter is true for the small- μ_i scenario, which is our major focus. Afterall, a mass eigenstate result may be considered as corresponding to taking an infinite summation over electroweak-state diagrams with all possible mass insertions. The comparison helps to illustrate better the physics hidden behind the complicated formulae. Besides a short letter by the present authors [3], some parts of the results here have also been discussed in Ref. [4].

A. Gluino Contributions

The most direct contributions come from a gaugino loop, as shown in Fig. 1. The diagram looks the same as the MSSM gluino and neutralino diagram with two gauge coupling vertices. As pointed out in our previous short letter [3], the new RPV contributions here is a simple result of the RPV LR squark mixings [*cf.* Eq.(8)]. In Ref. [3], we focused on the dominant gluino loop contribution. We first give some more details of that analysis, before going into the other contributions. Notice that though both the u and d quarks get EDM from gaugino loops in MSSM, only the d quark has the RPV contribution. The u -squark sector simply has no RPV LR mixings. In Fig. 1, as well as the subsequent diagrams, we use a family index

k for the external quark lines though only the $k = 1$ case would be the d (or u) quark EDM we are mainly concerned with. With the only possible exception of d_2 , which corresponds to the s quark [15], $k \neq 1$ results are not relevant for neutron EDM.

For illustrative purpose, we first neglect inter-family mixings among the squarks. The \tilde{d} mass-squared matrix, of the form given in Eq.(6) but reduced to one family, is hermitian and can be diagonalized by the unitary transformation

$$\mathcal{D}_d^\dagger \mathcal{M}_D^2 \mathcal{D}_d = \text{diag}\{M_{d-}^2, M_{d+}^2\}, \quad (16)$$

with

$$M_{d\mp}^2 = \frac{1}{2} \left[(\mathcal{M}_{LL}^2 + \mathcal{M}_{RR}^2) \mp \sqrt{(\mathcal{M}_{LL}^2 - \mathcal{M}_{RR}^2)^2 + 4|\mathcal{M}_{RL}|^2} \right], \quad (17)$$

and

$$\mathcal{D}_d = \begin{pmatrix} \cos \frac{\theta}{2} & -\sin \frac{\theta}{2} e^{-i\phi} \\ \sin \frac{\theta}{2} e^{-i\phi} & \cos \frac{\theta}{2} \end{pmatrix}. \quad (18)$$

Here, $\mathcal{M}_{RL}^2 = |\mathcal{M}_{RL}|^2 e^{i\phi}$ and the range of θ can be chosen so that $-\frac{\pi}{2} \leq \theta \leq \frac{\pi}{2}$ with $\tan\theta = \frac{2|\mathcal{M}_{RL}|^2}{\mathcal{M}_{LL}^2 - \mathcal{M}_{RR}^2}$. In terms of mass eigenstates \tilde{d}_- and \tilde{d}_+ , the gluino contribution to the EDM is then given by [16,20]

$$\left(\frac{d_d}{e} \right)_{\tilde{g}} = -\frac{2\alpha_s}{3\pi} \left[\text{Im}(\Gamma_d^{11}) \frac{M_{\tilde{g}}}{M_{d-}^2} \mathcal{Q}_{\tilde{d}} B\left(\frac{M_{\tilde{g}}^2}{M_{d-}^2}\right) + \text{Im}(\Gamma_d^{12}) \frac{M_{\tilde{g}}}{M_{d+}^2} \mathcal{Q}_{\tilde{d}} B\left(\frac{M_{\tilde{g}}^2}{M_{d+}^2}\right) \right] \quad (19)$$

where $\mathcal{Q}_{\tilde{d}} = -\frac{1}{3}$, $\Gamma_d^{1k} = \mathcal{D}_{d2k} \mathcal{D}_{d1k}^*$, giving $\text{Im}(\Gamma_d^{11}) = -\text{Im}(\Gamma_d^{12}) = \frac{1}{2} \sin\theta \sin\phi$, and

$$B(x) = \frac{1}{2(x-1)^2} \left[1 + x + \frac{2x \ln x}{(1-x)} \right]. \quad (20)$$

Introducing $M_d^2 = (M_{d-}^2 + M_{d+}^2)/2$ and using the identity relation $x B(x) = B(1/x)$, the gluino contribution becomes the often quoted

$$\left(\frac{d_d}{e} \right)_{\tilde{g}} = -\frac{2\alpha_s}{3\pi} \frac{M_{\tilde{g}}}{M_d^2} \mathcal{Q}_{\tilde{d}} \text{Im}(\delta_{\text{II}}^D) F\left(\frac{M_{\tilde{g}}^2}{M_d^2}\right), \quad (21)$$

where δ_{II}^D is \mathcal{M}_{RL}^2/M_d^2 (with \mathcal{M}_{RL}^2 restricted to the \tilde{d} family) and

$$F(x) = \frac{1}{(1-x)^3} \left[\frac{1+5x}{2} + \frac{(2+x)x \ln x}{(1-x)} \right]. \quad (22)$$

The EDM expression above is, in fact, the same as that of the MSSM case, except that δ_{II}^D , or equivalently \mathcal{M}_{RL}^2 , has now an extra RPV part. From the general result given in Eq.(8), we have for the \tilde{d} squark,

$$\delta_{11}^D M_d^2 = [A_d - \mu_0^* \tan\beta] m_d + \frac{\sqrt{2} M_W \cos\beta}{g_2} \delta A_{11}^D - \frac{\sqrt{2} M_W \sin\beta}{g_2} (\mu_i^* \chi_{i11}) . \quad (23)$$

Note that the $\mu_i^* \chi_{i11}$ term does contain nontrivial CP violating phases and gives RPV contribution to neutron EDM. Though the above analysis neglects inter-family mixings, it is clear from Fig. 1 that the LR squark mixing δ_{11}^D is what is directly responsible for the EDM contribution. Including inter-family mixings would complicate the mass eigenstate analysis but not modify the EDM result in any substantial way.

B. Neutralino-like Contributions

The contributions from the (electroweak) neutral gaugino loop, as shown in Fig. 1, are expected to be small, due to the much smaller gauge couplings. Apart from the neutral gaugino loops, there may be other neutralino-like loop contributions. In the MSSM, one has to consider possible contributions from the higgsino parts and the gaugino/higgsino mixing parts. Part of such contributions involve no LR squark mixing. The latter feature compensates for the smaller Yukawa type couplings involved [19]. Without R-parity, the gauginos and higgsinos mix with the leptons. We have seven neutral fermions, all among which give rise to a neutralino-like loop contribution.

The neutral fermion loop contribution to u and d quark EDM's is given by

$$\left(\frac{d_f}{e}\right)_{\chi^0} = \frac{\alpha_{\text{em}}}{4\pi \sin^2\theta_W} \mathcal{Q}_{\tilde{f}} \sum_{\tilde{f}\mp} \sum_{n=1}^7 \text{Im}(\mathcal{N}_{f n \mp}) \frac{M_{\chi_n^0}}{M_{\tilde{f}\mp}^2} B\left(\frac{M_{\chi_n^0}^2}{M_{\tilde{f}\mp}^2}\right) , \quad (24)$$

where

$$\begin{aligned} \mathcal{N}_{f n -} &= \left[-\sqrt{2} \{ \tan\theta_W (\mathcal{Q}_f - T_{3f}) \mathbf{X}_{1n} + T_{3f} \mathbf{X}_{2n} \} \mathcal{D}_{f11}^* - \frac{y_f}{g_2} \mathbf{X}_{bn} \mathcal{D}_{f21}^* - \delta_{b4} \frac{\chi_{k11}}{g_2} \mathbf{X}_{(k+4)n} \mathcal{D}_{d21}^* \right] \\ &\quad \cdot \left[\sqrt{2} \tan\theta_W \mathcal{Q}_f \mathbf{X}_{1n} \mathcal{D}_{f21} - \frac{y_f}{g_2} \mathbf{X}_{bn} \mathcal{D}_{f11} - \delta_{b4} \frac{\chi_{h11}}{g_2} \mathbf{X}_{(h+4)n} \mathcal{D}_{d11} \right] , \\ \mathcal{N}_{f n +} &= \left[-\sqrt{2} \{ \tan\theta_W (\mathcal{Q}_f - T_{3f}) \mathbf{X}_{1n} + T_{3f} \mathbf{X}_{2n} \} \mathcal{D}_{f12}^* - \frac{y_f}{g_2} \mathbf{X}_{bn} \mathcal{D}_{f22}^* - \delta_{b4} \frac{\chi_{k11}}{g_2} \mathbf{X}_{(k+4)n} \mathcal{D}_{d22}^* \right] , \\ &\quad \cdot \left[\sqrt{2} \tan\theta_W \mathcal{Q}_f \mathbf{X}_{1n} \mathcal{D}_{f22} - \frac{y_f}{g_2} \mathbf{X}_{bn} \mathcal{D}_{f12} - \delta_{b4} \frac{\chi_{h11}}{g_2} \mathbf{X}_{(h+4)n} \mathcal{D}_{d12} \right] ; \end{aligned} \quad (25)$$

with $b = 3 (4)$ for $T_{3f} = \frac{1}{2} (-\frac{1}{2})$ [*i.e.* for f being the u (d) quark], and y_f the corresponding Yukawa coupling [*cf.* Eq.(4)]. Recall, from Eq.(16), that \tilde{d}_{\mp} denotes the two \tilde{d} mass eigenstates and \mathcal{D}_d the diagonalizing matrix; likewise \tilde{u}_{\mp} and \mathcal{D}_u are the corresponding notations for the u quark case. Finally, the \mathbf{X}_{ij} 's are elements of the matrix that diagonalizes $\mathcal{M}_{\mathcal{N}}$, as defined explicitly in Appendix A. Note that the last term in each bracket of the $\mathcal{N}_{f n \mp}$ expressions is nonvanishing only for $f = d$, as indicated by the δ_{b4} symbol. Similar to the

case of the gaugino contributions, RPV contributions here exist only for the d quark EDM. We would like to emphasize that the formulae here represent the full 1-loop result for our generic supersymmetric SM.

Each expression for $\mathcal{N}_{fn\mp}$ above is a product of two terms, from the two loop vertices involving the L - and R -handed quark fields respectively, as a chirality flip has to be imposed within the loop. For each vertex, the three terms (in each bracket) correspond to the gauge, quark Yukawa, and RPV λ - couplings, with the last exist only for the d quark case. With two gauge coupling vertices, we have the gaugino diagram. Diagrams with a gauge and a y_f coupling is shown explicitly in Fig. 2 for the d quark, in which y_d is indicated by λ'_{0kk} , *i.e.* with the notation used in our superpotential and a general, unspecified, quark family index ($y_d \equiv \lambda'_{0ii}$ in our notation). Note that the diagrams require no LR mixing on the squark line. This is the important MSSM contribution, the chargino counterpart of which typically receives the major attention. The latter is numerically a bit larger. The type of diagrams has no RPV contribution.

There is, however, a RPV analog to Fig. 2. With the notation as given, this is obviously given by replacing λ'_{0kk} in the figure with a RPV λ'_{ikk} . This is shown explicitly in Fig. 3. From the latter figure, one can see the the first order result would come from a l_i^0 -gaugino mass mixing. However, under our formulation, the latter is vanishing [see Eq.(A8) in appendix A]. Looking at the type of contributions from our EDM formula here, while a λ'_{kii} coupling may not be small, it comes into the formula with a $\mathbf{X}_{(k+4)n}$. For $n = 1$ to 4, corresponding to the heavy neutralino mass eigenstates, $\mathbf{X}_{(k+4)n}$ measures a RPV mixing in the neutral fermion masses $\mathcal{M}_{\mathcal{N}}$. For $n = 5$ to 7, at least one of the three $\mathbf{X}_{(k+4)n}$'s is expected to be order one, but the corresponding (physical neutrino) mass eigenvalues $M_{\chi_n^0}$'s give a strong suppression factor in the resulted EDM contributions. There is also a further suppression from the mixing factor of the gaugino part, *e.g.* $\mathbf{X}_{1(k+4)}$. Furthermore, there is, potentially, a GIM cancellation among the seven mass eigenstates. Our numerical results, however, do show that the type of contributions is generally important. We will discuss the issue more carefully below, using the explicit example of its chargino-like counterpart.

Lastly, we come to the diagrams with no gauge vertices, which we show in Fig. 4 for the d quark, using our generic $\lambda'_{\alpha jk}$ notation. Such a diagram in the MSSM case are totally negligible due to the small Yukawa couplings involved and the suppression from the LR squark mixing required. In fact, as shown in the figure, the minimal mass insertion needed for the diagram corresponds to a Majorana mass among the l_α^0 's, which is again vanishing. The situation is similar to that of Fig. 3. The contribution is expected to be always dominated by the latter one.

C. Chargino-like Contributions

Next we come to the charged fermion counterpart. With similar notations as used in the neutral fermion loop formula above, the charged fermion loop contributions to u and d quark EDM can be written as

$$\left(\frac{d_f}{e}\right)_{\chi^-} = \frac{\alpha_{\text{em}}}{4\pi \sin^2\theta_W} \sum_{\tilde{f}'\mp} \sum_{n=1}^5 \text{Im}(\mathcal{C}_{f n \mp}) \frac{M_{\chi_n}^-}{M_{\tilde{f}'\mp}^2} \left[\mathcal{Q}_{\tilde{f}'} B\left(\frac{M_{\chi_n}^2}{M_{\tilde{f}'\mp}^2}\right) + (\mathcal{Q}_f - \mathcal{Q}_{\tilde{f}'}) A\left(\frac{M_{\chi_n}^2}{M_{\tilde{f}'\mp}^2}\right) \right], \quad (26)$$

for f being u (d) quark and f' being d (u), where

$$\begin{aligned} \mathcal{C}_{un-} &= \frac{y_u}{g_2} \mathbf{V}_{2n}^* \mathcal{D}_{d11} \left(-\mathbf{U}_{1n} \mathcal{D}_{d11}^* + \frac{y_d}{g_2} \mathbf{U}_{2n} \mathcal{D}_{d21}^* + \frac{\lambda'_{k11}}{g_2} \mathbf{U}_{(k+2)n} \mathcal{D}_{d21}^* \right), \\ \mathcal{C}_{un+} &= \frac{y_u}{g_2} \mathbf{V}_{2n}^* \mathcal{D}_{d12} \left(-\mathbf{U}_{1n} \mathcal{D}_{d12}^* + \frac{y_d}{g_2} \mathbf{U}_{2n} \mathcal{D}_{d22}^* + \frac{\lambda'_{k11}}{g_2} \mathbf{U}_{(k+2)n} \mathcal{D}_{d22}^* \right), \\ \mathcal{C}_{dn-} &= \left(\frac{y_d}{g_2} \mathbf{U}_{2n} + \frac{\lambda'_{k11}}{g_2} \mathbf{U}_{(k+2)n} \right) \mathcal{D}_{u11} \left(-\mathbf{V}_{1n}^* \mathcal{D}_{u11}^* + \frac{y_u}{g_2} \mathbf{V}_{2n}^* \mathcal{D}_{u21}^* \right), \\ \mathcal{C}_{dn+} &= \left(\frac{y_d}{g_2} \mathbf{U}_{2n} + \frac{\lambda'_{k11}}{g_2} \mathbf{U}_{(k+2)n} \right) \mathcal{D}_{u12} \left(-\mathbf{V}_{1n}^* \mathcal{D}_{u12}^* + \frac{y_u}{g_2} \mathbf{V}_{2n}^* \mathcal{D}_{u22}^* \right), \end{aligned} \quad (\text{only repeated index } i \text{ is to be summed}); \quad (27)$$

\mathbf{V} and \mathbf{U} are diagonalizing matrix of the charged fermion mass matrix \mathcal{M}_c as defined by Eq.(A2) in Appendix A; and the function $A(x)$ is given by

$$A(x) = \frac{1}{2(1-x)^2} \left(3 - x + \frac{2 \ln x}{1-x} \right). \quad (28)$$

The first part and the second part of Eq.(26) come from the cases in which the photon is emitted from the squark and fermion line inside the loop, respectively.

The basic feature of the $\mathcal{C}_{f n \mp}$'s terms is similar to the previous case of $\mathcal{N}_{f n \mp}$. They do not give rise to a charged gaugino loop contribution though, as R -handed quarks do not couple to \tilde{W}^\pm . Taking the available gauge coupling term within each $\mathcal{C}_{f n \mp}$ and a y_f to form an admissible product, we have the contribution corresponds to a diagram in Fig. 5. This is the dominating MSSM contribution, besides the equally competitive gluino one. Note again that no LR squark mixing is required the diagrams.

The diagram given in Fig. 5b has the RPV analog, which requires a $l_k^- - \tilde{W}^+$ mass insertion for the first order result, as shown in in Fig. 6. This is the $SU(2)$ partner of Fig. 3, something we promise to discuss in more detail. Looking at the contributions from our EDM formula here, we have again a λ'_{k11} coming in with a $\mathbf{U}_{(k+2)n}$, as verses the $\mathbf{X}_{(k+4)n}$ in the neutral case above. For $n = 1$ to 2 , corresponding to the heavy chargino mass eigenstates, $\mathbf{U}_{(k+2)n}$

measures a RPV mass mixing; it is order one for $n = k + 2$, but then $M_{\chi_n^-}$ is the small m_k , roughly the charged lepton mass. When one sums over for $n = 1$ to 5, it is easy to see that the result is proportional to the imaginary part of

$$\sum_{n=1}^5 \mathbf{V}_{1n}^* M_{\chi_n^-} \mathbf{U}_{(k+2)n} F_{BA}(M_{\chi_n^-}^2) \chi'_{k11} \quad (29)$$

where $F_{BA}(M_{\chi_n^-}^2)$ denotes a function on $M_{\chi_n^-}^2$ corresponding to the big bracket in Eq.(26) with functions B and A . If the F_{BA} could be factor out, together with χ'_{k11} , what is left inside the summation is nothing other than the $l_k^- \tilde{W}^+$ mass term, which is zero. This is a GIM-like cancellation mechanism, violated only to the extent that the loop integrals involved as given by the B and A functions are not universal. Our numerical calculations, however, show that the cancellation is very substantially violated, for generic values of chargino masses. Let us then look at the contribution from an individual mass eigenstate more closely. To get an analytical approximation, we used perturbative results for $\mathbf{U}_{(k+2)n}$ given in Appendix A. For the $n = 1$ and 2 parts in the above sum, we have

$$\mathbf{V}_{1a}^* \mu_k^* R_{R2a} F_{BA}(M_{ca}^2) \chi'_{k11} \quad (30)$$

with basically the same source of RPV complex phase as the gluino case, namely from $\text{Im}(\mu_k^* \chi'_{k11})$, except that it is one value of k for each diagram here. Note that the explicit M_{ca} factors are cancelled. The $n = k + 2$ part is largely suppressed due to a small mass eigenvalue, and the very small RPV R -handed mixing given by $\mathbf{V}_{1(k+2)}^*$ as shown in the appendix. Note that $\mathbf{V}_{1a}^* \simeq R_{R1a}^*$; and R_R is a 2×2 unitary matrix of order one elements. The expressions give an idea about the strength of the RPV contributions. In the limit of degenerate charginos, *i.e.* $M_{c1} = M_{c2}$, the F_{BA} function factors out of the sum of the $n = 1$ and 2 parts and the GIM-like cancellation is clearly illustrated, simply from the unitarity of R_R . In fact, our numerical calculations give a cancellation up to one part in 10^4 in such a situation if we enforce the condition. However, that requires a very substantial complex phase for μ_0 , hence of not most phenomenological interest. Another interesting point on the parameter space appears when the R_R matrix is diagonal, at which the contribution also goes away, as is obvious from the above expression. This happens at the point where the condition $M_2^* \sin\beta = \mu_0 \cos\beta$ is satisfied. The feature will be confirmed by our numerical results presented below.

The rest of the chargino-like contributions each has at least one y_f coupling and no gauge coupling vertex. Again, a LR squark mixing is required. However, the first order electroweak state diagrams, as depicted in Fig. 7, each has an admissible μ_α^* mass insertion on the fermion line. This looks a bit different from the neutral case above. Apart from an admissible RPV down-squark mixing for the case of the u quark, a RPV diagram again has the RPV parameter combination $\mu_k^* \chi'_{k11}$ (no sum). Except for large μ_i case, there is no

reason to expect the contribution to be of any significance compared to the other major contributions.

Our EDM formulae neglect inter-family mixings. However, general, unspecified, family indices are use in our figures to make the flavor structure transparent. We note that in the case of the contributions depicted by Fig. 4 and 7, off-diagonal $\mu_i^* \lambda_{ijk}$ type RPV mixings could play a role to let the higher family squarks mediating more important contributions to u and d quark EDM's.

D. The Quark-Scalar Loop Contributions in Brief

There are superpartners to the diagrams in Fig. 7, *i.e.* with quarks and charged scalars running inside the loop and a B_α coupling on the scalar line. These diagrams are shown explicitly in Fig. 8. Here, the $\alpha = 0$ case gives the MSSM charged Higgs contribution. With the B_i 's, we have the RPV contributions. Analytically, we have the formula

$$\left(\frac{d_f}{e}\right)_{\phi^-} = \frac{\alpha_{\text{em}}}{4\pi \sin^2\theta_W} \sum'_m \sum_{n=1}^3 \text{Im}(\mathcal{C}_{inm}^L \mathcal{C}_{inm}^{R*})_f \frac{M_{f'n}}{M_{\tilde{\ell}_m}^2} \left[(\mathcal{Q}_f - \mathcal{Q}_{f'}) B\left(\frac{M_{f'n}^2}{M_{\tilde{\ell}_m}^2}\right) + \mathcal{Q}_{f'} A\left(\frac{M_{f'n}^2}{M_{\tilde{\ell}_m}^2}\right) \right], \quad (31)$$

where for $f = u$

$$\begin{aligned} \mathcal{C}_{inm}^R &= \frac{y_d}{g_2} \mathcal{D}_{1m}^{l*} + \frac{\lambda_{i1k}^*}{g_2} \mathcal{D}_{(i+2)m}^{l*}, \\ \mathcal{C}_{inm}^L &= \frac{y_u}{g_2} \mathcal{D}_{2m}^{l*}; \end{aligned} \quad (32)$$

and for $f = d$

$$\begin{aligned} \mathcal{C}_{inm}^R &= \frac{y_u}{g_2} \mathcal{D}_{1m}^l, \\ \mathcal{C}_{inm}^L &= \frac{y_d}{g_2} \mathcal{D}_{2m}^l + \frac{\lambda_{ij1}}{g_2} \mathcal{D}_{(i+2)m}^l; \end{aligned} \quad (33)$$

with the \sum'_m denotes sum over (seven) non-zero mass eigenstates of the charged scalar, *i.e.* the unphysical Goldstone mode is dropped from the sum; \mathcal{D}^l being the diagonalization matrix, *i.e.* $\mathcal{D}^{l\dagger} \mathcal{M}_E^2 \mathcal{D}^l = \text{diag}\{M_{\tilde{\ell}_m}^2, m = 1 \text{ to } 8\}$.

With similar notations as used in the charged scalar loop formula above, the neutral scalar loop contributions to the EDM can be written as

$$\left(\frac{d_f}{e}\right)_{\phi^0} = \frac{\alpha_{\text{em}}}{4\pi \sin^2\theta_W} \sum'_m \sum_{n=1}^3 \text{Im}(\mathcal{N}_{inm}^L \mathcal{N}_{inm}^{R*}) \frac{M_{fn}}{M_{S_m}^2} \mathcal{Q}_f A\left(\frac{M_{fn}^2}{M_{S_m}^2}\right), \quad (34)$$

where for $f = u$

$$\begin{aligned}
\mathcal{N}_{inm}^R &= -\frac{y_u}{g_2} \frac{1}{\sqrt{2}} [\mathcal{D}_{1m}^s - i \mathcal{D}_{6m}^s], \\
\mathcal{N}_{inm}^L &= -\frac{y_u}{g_2} \frac{1}{\sqrt{2}} [\mathcal{D}_{1m}^s + i \mathcal{D}_{6m}^s];
\end{aligned} \tag{35}$$

and for $f = d$

$$\begin{aligned}
\mathcal{N}_{inm}^R &= -\frac{y_d}{g_2} \frac{1}{\sqrt{2}} [\mathcal{D}_{2m}^s - i \mathcal{D}_{7m}^s] - \frac{\lambda_{h1k}^*}{g_2} \frac{1}{\sqrt{2}} [\mathcal{D}_{(h+2)m}^s - i \mathcal{D}_{(h+7)m}^s], \\
\mathcal{N}_{inm}^L &= -\frac{y_d}{g_2} \frac{1}{\sqrt{2}} [\mathcal{D}_{(i+2)m}^s + i \mathcal{D}_{(i+7)m}^s] - \frac{\lambda_{hk1}}{g_2} \frac{1}{\sqrt{2}} [\mathcal{D}_{(h+2)m}^s + i \mathcal{D}_{(h+7)m}^s];
\end{aligned} \tag{36}$$

with again the unphysical Goldstone mode to be dropped from the sum over scalar mass eigenstates (nine non-zero); and the obvious notation $(\mathcal{D}^s)^T \mathcal{M}_s^2 \mathcal{D}^s = \text{diag}\{M_{s_m}^2, m = 1 \text{ to } 10\}$, \mathcal{D}^s being an orthogonal matrix.

Note that the formulae given in this section, like those in the rest of the paper, have neglected CKM mixings among the quarks. However, unlike the previous cases, the mixings play an important role in the quark-scalar loop contributions. This is a result of the fact that the EDM contributions have a fermion, quark in this case, mass dependence, and the mass hierarchy among the quarks. For instance, the top loop is expected to give a dominating contribution. Generalizing the formulae to include the mixings is straight forward. The formulae as they are given in this sub-section do give a basic illustration of the major analytical structure of the type of contributions. We will only discuss the features briefly here.

Diagrams depicting the neutral scalar contributions are like superpartners of the type of diagram given in Fig. 3. This is shown in Fig. 8, for both the u and d quark cases. Such a diagram could be interpreted as requiring a Majorana like scalar mass insertion. The latter is first considered in Ref. [11], under the name of Majorana masses from the “sneutrinos”. However, looking at it from the present framework, these diagrams should be interpreted as having Majorana like mass insertions among the l_α^0 's (and h_u^0 for the case of u quark).

The quark-scalar loop given by the above formulae are not quite considered in the case of MSSM, as they would be suppressed by the very small Yukawa couplings. The RPV analog, however, admits much more general flavor structure, as illustrated in the diagrams in Fig. 8 and 9. For instance, a top quark with order one Yukawa coupling could be contributing to the d quark EDM through a charged scalar loop. Contributions discussed in this sub-section depend on a much larger number of parameters, through the scalar mass matrices [*cf.* Eqs.(9)-(15)]. The extra parameters are related directly to Higgs physics. The important RPV B_i parameters here have strong connection with the μ_i 's [23]. Furthermore, the more general flavor structure involved means that the EDM contributions involve λ_{ijk} couplings that would also have important roles to play in the related processes of $b \rightarrow s \gamma$ and $b \rightarrow d \gamma$. We will leave all these issues for later studies.

V. RESULT FROM NUMERICAL CALCULATIONS

We now discuss the results we obtained by a careful numerical implementation of our EDM formulae discussed above, with explicit numerical diagonalization of all the mass matrices involved. Note that the quark-scalar loop contribution is not included in the numerical study. Hence, we are concentrating on the SUSY contributions in the presence of RPV couplings, to be compared directly with the MSSM results.

As discussed in the previous section, the imaginary part of the combination $\mu_i^* \chi_{i11}$ is what is central to the RPV 1-loop EDM contributions. To simplify the discussion, we single out μ_3 and χ_{311} and put the corresponding parameters for $i = 1$ and 2 to be essentially zero³. This applies to all the numerical results discussed here.

We give in Table 1 some detailed numerical results for four illustrative sample cases. To first focus on the RPV contributions which we are particularly interested in, cases A and B in the table have all complex phases in the R-parity conserving part switched off. For the RPV parameters μ_3 and χ_{311} , we take the former being real and put a $\pi/4$ complex phase into the latter. However, we want to emphasize again that only the phase of the combination $\mu_3^* \chi_{311}$ matters. For instance, we have checked explicitly that shifting the overall phase, or a part of it into μ_3 instead gives exactly the same results. The difference between case A and case B is only in the value of μ_3 chosen. Case A has small μ_3 , at 10^{-3} GeV. This is the small- μ_i scenario, corresponding to a sub-eV mass for ν_τ as suggested, but far from mandated, by the result from the Super-Kamiokande (super-K) experiment [24]⁴. In contrast, case B has $\mu_3 = 1$ GeV. Note that imposing the 18.2 MeV experimental bound [25] for the mass of ν_τ still admits a relatively large μ_3 , especially for a large $\tan\beta$. Reading from the results in Ref. [7], the bound is ~ 7 GeV at $\tan\beta = 2$ and ~ 300 GeV at $\tan\beta = 45$. As for χ_{311} , the best bound on the (from $\tau \rightarrow \pi\nu$) is around $0.05 \sim 0.1$ [26]. Hence, case B is still within admissible region of RPV parameter space, though beyond the more popular small- μ_i scenario. Moreover, we have illustrated in Ref. [3] that the χ_{311} bound is not strengthened even when the above mentioned stringently interpreted neutrino mass bound from super-K is imposed on the 1-loop neutrino mass contribution obtainable from the parameter. We, therefore, use the same χ_{311} magnitude of 0.05 for both, and actually all, cases in the table.

Apart from the overall EDM results, we are interested in understanding the relative values

³In the actual calculation, tiny but nonzero values are used to avoid possible problems of numerical manipulations.

⁴Note that $\mu_3 \cos\beta \sim 10^{-4}$ GeV is actually expected, allowing a larger μ_3 in the case of large $\tan\beta$ (see, for example, Ref. [8]). Our choice of μ_3 value here is, admitted, quite pushing the limit. It is mainly for illustrative purpose.

of the different pieces of the contributions discussed in the previous section. Hence, we show, in the table, the contribution from individual pieces, identified by the couplings of the two loop vertices given inside the first column. They can be matched easily with terms in the formulae. The corresponding Feynman diagrams are marked whenever they are available, inside the next column. Column three of the table indicates whether a LR squark mixing is involved in the contribution, and if the RPV part of the latter is available and needed for a RPV contribution to come in. The individual contribution that is vanishing under cases A and B corresponds to a diagram which does not admit a RPV contribution. This is the case for most entries for the u quark EDM and the entries for the chargino-like loop contributions to the d quark EDM not involving a λ coupling loop vertex. Note that a chargino-like loop for the d quark involves LR mixing of \tilde{u} which has no RPV part. As emphasized in Ref. [3], the RPV contributions are far more prominent for the d quark part. Again, λ_{311} is the only nonzero λ coupling introduced, though we use more general λ notations within the table to remind the readers of the more general flavor structure admissible, as illustrated in the previous section. We also recall from the discussions there that a λ'_{ijk} coupling always comes into the EDM picture accompanied by a RPV fermion mixing matrix element proportional to μ_i^* . The EDM results are basically in direct proportion to $\mu_i^* \lambda_{i11}$ over the whole region of parameter space studied. Hence, comparing cases A and B, we see an *exact* factor of 10^{-3} suppression in case A for contributions with one λ loop vertex, or involving a RPV LR mixing, except the very small y_d^2 term.

The overall EDM numbers of case A is still below the present experimental bound. This is smaller than a naive estimated result, as given for the gluino contribution by Eq.(21) for example, due to some unavoidable partial cancellation. For the gluino result, in particular, the cancellation is between the contributions from the two squark mass eigenstates, which would in fact be exact when the latter states are degenerate. This slightly weakens the EDM bound of the previously estimated $\text{Im}(\mu_i^* \lambda_{i11}) < 10^{-6} \text{ GeV}$ result given in Ref. [3]. However, we see here that with unification type relationship imposed on the gaugino masses, we are likely to have a chargino contribution larger than the gluino one. Moreover, the values of the other SUSY parameters chosen here can be pushed to increase the EDM number, as discussed below. As for case B, the numbers are beyond the experimental bound, indicating that the RPV parameters, or the involved complex phase, would have to be further constrained.

Case C of Table 1 gives a different scenario. Here, we allow complex phases, and hence EDM contributions, of the MSSM part. The RPV parameters are set real. This ensures no RPV EDM contribution from diagrams involving LR squark mixings. The interesting point to note here is that there is indeed nonzero RPV contributions from diagrams with RPV loop vertices. In particular, the RPV chargino contribution to d quark EDM is comparable in magnitude to its R-parity conserving counterpart. Hence, the existence of nonzero RPV parameters, even real, would change the EDM story of the MSSM. In other words, in the

presence of a complex μ_0 , the neutron EDM bound actually constrains the magnitude of the the combination of RPV parameters given by $\mu_i^* \lambda'_{i\text{II}}$, instead of just the phase of it. This is an important feature that has not been pointed out before. Result of the case serves, otherwise, as a checking against other neutron EDM studies of MSSM.

Finally, we illustrate in case D a specific situation with both R-parity conserving and violating phases. Here, we pick a case of large $\tan\beta$, for which the larger μ_3 is still within the limit admitted by the stringent interpretation of the super-K bound [8]. The negative phase chosen for $\lambda'_{3\text{II}}$ gives a cancellation between the two contributions to the imaginary part of LR d squark mixing. This directly resulted in suppression of the EDM contributions involving the latter, such as the d quark gluino loop. The very interesting point to note here is that the same cancellation is resulted in the chargino and neutralino loop ⁵, as explicitly illustrated by the entries of the terms with loop vertex couplings given by $g \cdot y_d$ and $g \cdot \lambda'_{i\text{II}}$ in both cases. Essentially, when we have a small imaginary part for $\mu_\alpha^* \lambda'_{\alpha\text{II}}$ (recall: $\alpha = 0$ to 3 , $\lambda'_{0\text{II}} \equiv y_d$), *e.g.* an internal cancellation among the summed terms, the corresponding d quark EDM contributions, involving the suppressed LR mixing [*cf.* first line of Eq.(8)] or otherwise, are all suppressed. The chargino and neutralino contributions actually reflex some proportionality to the (F -term) LR mixings in a way similar to the gluino contribution. Note however that the A -term phase contributes to the gluino diagram while leaving the chargino-like and neutralino-like diagram not quite affected, hence spoils the simultaneous cancellation achieved in Case D results here. Besides, as the u quark EDM results do not have the complete matching analog between the R-parity conserving and violating contributions; the simultaneous cancellation mechanism does not exist there. In the current case, the u quark parts are suppressed due to large $\tan\beta$.

The above sample cases illustrate some of the interesting features about the RPV contributions to neutron EDM. Next, we discuss how the EDM result varies with some basic parameters. Fig. 10 shows a logarithmic plot of (the magnitude of) the RPV neutron EDM result for μ_0 value between ± 2000 GeV, with the other parameters set at the same values as Case A in Table 1. Note the the small $|\mu_0|$ ($\lesssim 70$ GeV) region has been ruled out for having too light a chargino mass [7]. As noted in the previous section, the gluino and chargino contributions compete with one another with one of them being dominating in a region of the parameter space. The dip on the C -line for the chargino-like contribution corresponds

⁵A chargino (neutralino) contribution means exactly the one with a physical chargino (neutralino) running in the loop. Hence, it is only a part of the fermion mass eigenstate sum in the analytical formulae given for the (total) chargino-like and neutralino-like loop contributions. The latter, as illustrated in the previus section, are generally dorminated by the chargino and neutralino mass eigenstates.

to the case of vanishing R -handed mixing among the charginos, *i.e.* when the condition $M_2 \sin\beta + \mu_0 \cos\beta = 0$ is satisfied, as noted above, where the contribution essentially vanishes. Note that in the region to the left of the point, the contribution becomes negative. The G -line is horizontal, as the gluino diagram result has no dependence on μ_0 .

Next, we check the $\tan\beta$ dependence and compare with the MSSM result. In Fig. 11, we give again a plot of the (total) EDM result for the same set of input parameters as in Case A of Table 1, while varying $\tan\beta$. This is the line marked as “RPV only”. The numerical result confirms our earlier discussion that there is not much sensitivity in $\tan\beta$ [*cf.* expression (30)], as indicated by the flatness of the line. This is in good contrast to the MSSM result, also based on the same set of input parameter except with nonvanishing phases for A and μ_0 . The third line, marked by “GSSM”, gives the total result obtained from our formulae given above with the same nonvanishing phases for A and μ_0 , as well as χ_{3H} . Note that the GSSM result here is more than the sum of the two other lines, due to the presence of RPV contribution even in the limit of a real χ_{3H} , as discussed for Case C of Table 1 above. The $\tan\beta$ dependence, or the lack of, illustrated in the figure is quite generic, in a wide region of the parameter space.

The dip on the GSSM line of Fig. 11 represents a point where the overall contribution is small as a result of the cancellation among the different pieces, a feature that attracts lots of attention lately in the case of MSSM (see, for example, Ref. [27] and Refs. [20,21]). The MSSM line in the figure has the dip only beyond the range of $\tan\beta$ value shown. As also pointed out in a little bit different setting (with vanishing A phase) for Case D of Table 1 above, the new RPV contribution modifies the overall picture and provides possible new cancellation mechanism. The cancellation feature is better illustrated in Fig. 12, in which variation against the A -term phase is shown, for a small and a large $\tan\beta$ cases. Again our GSSM result is compared with that of MSSM. One can see that the presence of the RPV contribution shifts the cancellation points substantially. Finally, we show variations of the result as a function of the gaugino mass parameters here represented by M_2 . This is given in Fig. 13 with the four lines marked and correspond to the four cases of Fig. 12, each with the A -term phase set at the position of the dip as given in the latter figure.

VI. CONCLUSION

We have given explicit formulae and detailed discussions on the full 1-loop contributions to quark EDM for the generic supersymmetric SM (without R-parity). The extra, RPV, contributions are interesting additions to the R-parity conserving part. We have given results from an exact numerical study, illustrating various novel aspects. Our formulation emphasizes the universal structure of the R-parity conserving and violating parts, which also illustrates itself well in the results. The experimental bound on the neutron EDM is hence

established as an important source of constraints on the model parameter space, including the RPV part. The 1-loop RPV contributions always involves the particular combination of RPV parameters given by $\mu_i^* \lambda_{i\text{II}}$, with little sensitivity to the value of $\tan\beta$. So far as the RPV parameters are concerned, this combination is well constrained by the EDM bound. This applies not only to the complex phase, or imaginary part of, the combination. Real $\mu_i^* \lambda_{i\text{II}}$ contribute in the presence of complex phases in the chargino and neutralino mass entries. Studies with either the μ_i - or the λ' - type couplings assumed to be zero miss the class of very interesting phenomenological features of SUSY without R-parity.

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APPENDIX A: NOTES ON THE FERMION MASSES

Under the SVP, the (color-singlet) charged fermion mass matrix is given by the simple form :

$$\mathcal{M}_c = \begin{pmatrix} M_2 & \sqrt{2} M_W \cos\beta & 0 & 0 & 0 \\ \sqrt{2} M_W \sin\beta & \mu_0 & \mu_1 & \mu_2 & \mu_3 \\ 0 & 0 & m_1 & 0 & 0 \\ 0 & 0 & 0 & m_2 & 0 \\ 0 & 0 & 0 & 0 & m_3 \end{pmatrix}, \quad (\text{A1})$$

with explicit bases for right-handed and left-handed states given by $(-i\tilde{W}^+, \tilde{h}_u^+, l_1^+, l_2^+, l_3^+)$ and $(-i\tilde{W}^-, l_0^-, l_1^-, l_2^-, l_3^-)$, respectively. Here, we allow M_2 and all four μ_α parameters to be complex, though the m_i 's are restricted to be real, for reason that would become clear below. Obviously, each μ_i parameter here characterizes directly the deviation of the l_i^- from the corresponding physical charged lepton ($\ell_i = e, \mu$, and τ), *i.e.* light mass eigenstates. For any set of other parameter inputs, the m_i 's can then be determined, through a simple numerical procedure, to guarantee that the correct mass eigenvalues of m_e, m_μ , and m_τ are obtained — an issue first addressed and solved in Ref. [7]. The latter issue is especially important when μ_i 's not substantially smaller than μ_0 are considered. Such an odd scenario is not definitely ruled out [7]. However, for the more popular of small- μ_i scenario, we have $l_i^- \approx \ell_i^-$, and deviations on m_i 's from the (real) ℓ_i masses are negligible. Note that deviation of l_i^+ from ℓ_i^+ is quite negligible in any case.

We introduce unitary matrices \mathbf{V} and \mathbf{U} diagonalizing the R - and L -handed states with

$$\mathbf{V}^\dagger \mathcal{M}_c \mathbf{U} = \text{diag}\{M_{\chi_n}^-\} \equiv \text{diag}\{M_{c1}, M_{c2}, m_e, m_\mu, m_\tau\}. \quad (\text{A2})$$

Here, the mass eigenvalues $M_{\chi_n}^-$ with $n = 1$ and 2 , *i.e.* M_{c1} and M_{c2} , are the chargino masses. Consider further

$$R_R^\dagger \begin{pmatrix} M_2 & \sqrt{2} M_W \cos\beta \\ \sqrt{2} M_W \sin\beta & \mu_0 \end{pmatrix} R_L = \text{diag}\{M_{c1}^o, M_{c2}^o\} \quad (\text{A3})$$

with M_{c1}^o and M_{c2}^o being the chargino masses in the $\mu_i = 0$ limit. One can then write the diagonalizing matrices in the block form

$$\mathbf{V} = \begin{pmatrix} R_R & -R_R V^\dagger \\ V & I_{3 \times 3} \end{pmatrix} \quad \text{and} \quad \mathbf{U} = \begin{pmatrix} R_L & -R_L U^\dagger \\ U & I_{3 \times 3} \end{pmatrix}. \quad (\text{A4})$$

For $\mu_i \ll M_{ca}^o$ ($a = 1$ and 2), a block perturbative diagonalization can be performed directly on the matrix \mathcal{M}_c to obtain the following simple result :

$$\begin{aligned}
U_{(i+2)1} &\simeq \frac{\mu_i^*}{M_{c1}} R_{R21} , \\
U_{(i+2)2} &\simeq \frac{\mu_i^*}{M_{c2}} R_{R22} , \\
V_{(i+2)a} &\simeq \frac{m_i}{M_{ca}} U_{(i+2)a} \quad (a = 1 \text{ and } 2) .
\end{aligned} \tag{A5}$$

Elements in the R_R and R_L matrices are all expected to be of order one. The above expressions illustrate the magnitudes of given matrix elements involving RPV mixings, as well as their dependence on the RPV parameters. We note that the L -handed mixings are roughly measured by the ratio of a μ_i to the chargino mass scale, while R -handed mixings are further suppressed by a charged lepton to chargino mass ratio, hence quite negligible under most consideration.

Note that notation here is different from those given in Ref. [7], and many others in the literature. More explicitly, we have R - and L -handed mass eigenstates given by $(\chi_{+n}) = \mathbf{V}^T [-i\tilde{W}^+, \tilde{h}_u^+, l_1^+, l_2^+, l_3^+]^T$ and $(\chi_{-n}) = \mathbf{U}^\dagger [-i\tilde{W}^-, l_0^-, l_1^-, l_2^-, l_3^-]^T$; which form the five Dirac fermions $\chi_n^- = \begin{pmatrix} \chi_{-n} \\ \chi_{+n}^\dagger \end{pmatrix}$. The $U_{(i+2)a}$ elements as given above show no obvious dependence on $\tan\beta$, though some nontrivial dependence is expected through the R_{R2a} elements. Our exact numerical result also indicates a weak sensitivity on the $\tan\beta$ value here. On the other hand, we have, from Eqs.(A4) and (A5),

$$\begin{aligned}
U_{a(i+2)} &= -\mu_i \cdot \left[R_L (\text{diag}\{M_{c1}^o, M_{c2}^o\})^{-1} R_R^\dagger \right]_{a2} , \\
V_{a(i+2)} &= -\mu_i \cdot \left[R_R (\text{diag}\{M_{c1}^o, M_{c2}^o\})^{-2} R_R^\dagger \right]_{a2} ;
\end{aligned} \tag{A6}$$

giving the result

$$\begin{aligned}
U_{1(i+2)} &= \frac{\mu_i \sqrt{2} M_W \cos\beta}{M_0^2} , \\
U_{2(i+2)} &= -\frac{\mu_i M_2}{M_0^2} , \\
V_{1(i+2)} &= \mu_i m_i \frac{\sqrt{2} M_W (M_2^* \sin\beta + \mu_0 \cos\beta)}{|M_0|^4} , \\
V_{2(i+2)} &= -\mu_i m_i \frac{(|M_2|^2 + 2 M_W^2 \cos^2\beta)}{|M_0|^4} ,
\end{aligned} \tag{A7}$$

where

$$M_0^2 \equiv \mu_0 M_2 - M_W^2 \sin 2\beta = M_{c1}^o M_{c2}^o .$$

These RPV elements correspond to those given in Ref. [7] (with all complex phases neglected), where the crucial $\cos\beta$ dependence of the non-standard Z^0 -boson couplings of the

physical charged leptons ($\ell_i \equiv \chi_{i+2}$) through $\mathbf{U}_{1(i+2)}$ is emphasized. The difference between $\mathbf{U}_{(i+2)a}$ and $\mathbf{U}_{a(i+2)}$ is hence very important.

The 7×7 Majorana mass matrix for the neutral fermion can be written as

$$\mathcal{M}_{\mathcal{N}} = \begin{pmatrix} M_1 & 0 & M_Z \sin\theta_W \sin\beta & -M_Z \sin\theta_W \cos\beta & 0 & 0 & 0 \\ 0 & M_2 & -M_Z \cos\theta_W \sin\beta & M_Z \cos\theta_W \cos\beta & 0 & 0 & 0 \\ M_Z \sin\theta_W \sin\beta & -M_Z \cos\theta_W \sin\beta & 0 & -\mu_0 & -\mu_1 & -\mu_2 & -\mu_3 \\ -M_Z \sin\theta_W \cos\beta & M_Z \cos\theta_W \cos\beta & -\mu_0 & 0 & 0 & 0 & 0 \\ 0 & 0 & -\mu_1 & 0 & 0 & 0 & 0 \\ 0 & 0 & -\mu_2 & 0 & 0 & 0 & 0 \\ 0 & 0 & -\mu_3 & 0 & 0 & 0 & 0 \end{pmatrix}, \quad (\text{A8})$$

with explicit basis $(-i\tilde{B}, -i\tilde{W}, \tilde{h}_u^0, \tilde{h}_d^0, l_1^0, l_2^0, l_3^0)$. Note that $\tilde{h}_d^0 \equiv l_0^0$; and, from the above discussion of the charged fermions, we have, for small μ_i 's, $(l_1^0, l_2^0, l_3^0) \approx (\nu_e, \nu_\mu, \nu_\tau)$. The symmetric, but generally non-hermitian, matrix can be diagonalized by using unitary matrix \mathbf{X} such that

$$\mathbf{X}^T \mathcal{M}_{\mathcal{N}} \mathbf{X} = \text{diag}\{M_{\chi_n^0}\}. \quad (\text{A9})$$

Again, the first part of the mass eigenvalues, $M_{\chi_n^0}$ for $n = 1$ to 4 here, gives the heavy states, *i.e.* neutralinos. The last part, $M_{\chi_n^0}$ for $n = 5$ to 7 are hence physical neutrino masses at tree-level.

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Table caption :

Table 1 — Numerical 1-loop neutron EDM results from SUSY without R-parity, for four illustrative cases. All EDM numbers are in e cm. Note that the quark EDM numbers are direct output from the numerical program applying our quark dipole formulae; while the neutron EDM numbers are from the valence quark model formula, as given in Eq.(1). R-parity violating parameters not given are taken as essentially zero. All parameters are taken real except those with complex phases explicitly listed in each case, where the real number(s) listed then give the magnitude(s). Parameter A here means a common A_u and A_d . Only M_2 is shown for the gaugino masses, the other are fixed by the unification relationship. Explicitly, we use $M_1 = 0.5 M_2$ and $M_3 = 3.5 M_2$. The first column under “EDM Results” gives couplings of the the loop vertices involved. A ‘ g ’ indicates either one of the electroweak gauge coupling, while a λ coupling means one with the appropriate admissible flavor indices. In the explicit results of the four cases, the latter is always a λ_{311} . The second column gives the reference Feynman diagram figures, when available. The third column indicates whether the particular contribution involves a LR squark mixing. In the case that the mixing is involved and a R-parity violating (RPV) one is involved in generating a RPV EDM contribution, it is marked with ‘RPV’.

Figure captions :

Fig. 1 — Diagram for d quark EDM from gaugino loop.

Fig. 2 — Diagrams with neutral gaugino/higgsino mixing for d quark EDM.

Fig. 3 — R-parity violating neutralino-like loop diagrams for d quark EDM. Naive electroweak state analysis suggests that such a diagram is proportional to a vanishing l_k^0 -gaugino mass mixing.

Fig. 4 — Diagram for d quark EDM suggesting involvement of Majorana masses among the l_α^0 or “neutrinos”.

Fig. 5 — Diagrams for u and d quark EDM’s with charged gaugino/higgsino mixing.

Fig. 6 — R-parity violating chargino-like loop diagram for u and d quark EDM’s. Naive electroweak state analysis suggests that the diagrams are proportional to the vanishing $l_k^- - \tilde{W}^+$ mass term.

Fig. 7 — Diagrams for u and d quark EDM’s with a μ_α mass insertion.

Fig. 8 — Diagrams for u and d quark EDM’s with a B_α scalar mass insertion.

Fig. 9 — Diagrams with a Majorana-like scalar mass insertion for u and d quark EDM’s.

For the u quark case, the direct Majorana-like h_u mass insertion is explicitly shown. For the d quark case, the corresponding direct h_d mass insertion is obvious, for $\alpha = \beta = 0$; for α and/or β nonzero, the naive direct result from the diagram would vanish, due to the vanishing VEV's.

Fig. 10 — Logarithmic plot of (the magnitude of) the RPV neutron EDM result for μ_0 value between ± 2000 GeV, with the other parameters set at the same values as Case A in Table 1. The lines marked by G , C , N , and “Total” gives the complete gluino, chargino-like, neutralino-like, and total (*i.e.* sum of the three) contributions, respectively. Note that the values of the N contributions and those of the C line for $\mu_0 < -900$ GeV are negative.

Fig. 11 — Logarithmic plot of (the magnitude of) the neutron EDM result verses $\tan\beta$. We show here the MSSM result, our general result with RPV phase only, and the generic result with complex phases of both kinds. In particular, the A and μ_0 phases are chosen as 7° and 0.1° respectively, for the MSSM line. They are zero for the RPV-only line, with which we have a phase of $\frac{\pi}{4}$ for χ_{3H} . All the given nonzero values are used for the three phases for the generic result (from our complete formulae) marked by GSSM. Again, the other unspecified input parameters are the same as for Case A of Table 1.

Fig. 12 — Logarithmic plot of (the magnitude of) the neutron EDM result verses θ_A (the complex phase for the A parameter). The four line shown are characterized by the $\tan\beta$ values (3 or 50) used and whether it is for our GSSM result (again with a phase of $\frac{\pi}{4}$ for χ_{3H}) — marked by G , or the result for MSSM — marked by M . Again, the χ_{3H} phase is set at $\frac{\pi}{4}$ for the G lines, and the μ_0 phase at 0.1° for all; the other unspecified input parameters are the same as for Case A of Table 1.

Fig. 13 — Logarithmic plot of (the magnitude of) the neutron EDM result verses M_2 . The four lines correspond to the four cases of Fig. 12 each with θ_A set at the dip location, *i.e.* $G-3$ for GSSM at $\tan\beta = 3$ with $\theta_A = 2^\circ$, $M-3$ for MSSM at $\tan\beta = 3$ with $\theta_A = -1^\circ$, $G-50$ for GSSM at $\tan\beta = 50$ with $\theta_A = 20^\circ$, $M-50$ for MSSM at $\tan\beta = 50$ with $\theta_A = 3^\circ$. All other unspecified input parameters are the same as for Fig. 12.

<u>Choice of Parameters</u>						
$\tilde{m}_Q = 300 \text{ GeV}, \tilde{m}_u = \tilde{m}_d = 200 \text{ GeV}, A = M_2 = 300 \text{ GeV}, \mu_0 = -300 \text{ GeV}$						
$\tan\beta$		3	3	3	50	
μ_3		$1 \cdot 10^{-3} \text{ GeV}$	1 GeV	1 GeV	$5 \cdot 10^{-3} \text{ GeV}$	
χ_{311}		0.05	0.05	0.05	0.05	
(complex phases)		$\chi_{311}(\pi/4)$	$\chi_{311}(\pi/4)$	$\mu_0(0.5^\circ), A(10^\circ)$	$\mu_0(0.02^\circ), \mu_3(-\pi/4)$	
<u>EDM RESULTS :-</u>						
couplings	Fig.	LR -mixing	Case A	Case B	Case C	Case D
<u>d quark EDM</u>						
Gluino loop : -						
α_s	1	RPV	$8.8 \cdot 10^{-28}$	$8.8 \cdot 10^{-25}$	$-3.9 \cdot 10^{-26}$	$-6.7 \cdot 10^{-29}$
Neutralino-like loop : -						
g^2	1	RPV	$-1.9 \cdot 10^{-29}$	$-1.9 \cdot 10^{-26}$	$8.3 \cdot 10^{-28}$	$2.7 \cdot 10^{-30}$
$g \cdot y_d$	2	no	~ 0	~ 0	$-1.6 \cdot 10^{-27}$	$-1.2 \cdot 10^{-27}$
$g \cdot \chi_{i11}$	3	no	$-1.0 \cdot 10^{-28}$	$-1.0 \cdot 10^{-25}$	$1.1 \cdot 10^{-27}$	$1.1 \cdot 10^{-27}$
y_d^2	4	RPV	$9.7 \cdot 10^{-37}$	$9.7 \cdot 10^{-34}$	$-3.9 \cdot 10^{-35}$	$-2.6 \cdot 10^{-33}$
$y_d \cdot \chi'_{ijk}$	4	yes	$-1.7 \cdot 10^{-36}$	$-1.7 \cdot 10^{-33}$	$8.5 \cdot 10^{-35}$	$2.5 \cdot 10^{-33}$
two χ'_{ijk}	4	yes	$-2.1 \cdot 10^{-39}$	$-3.4 \cdot 10^{-34}$	$9.0 \cdot 10^{-35}$	$-8.6 \cdot 10^{-37}$
Chargino-like loop : -						
$g \cdot y_d$	5	no	~ 0	0	$2.5 \cdot 10^{-26}$	$1.7 \cdot 10^{-26}$
$g \cdot \chi'_{i11}$	6	no	$2.1 \cdot 10^{-27}$	$2.1 \cdot 10^{-24}$	$-1.3 \cdot 10^{-26}$	$-1.7 \cdot 10^{-26}$
$y_u \cdot y_d$	7	yes	~ 0	0	$-2.7 \cdot 10^{-34}$	$-8.0 \cdot 10^{-36}$
$y_u \cdot \chi'_{ijk}$	7	yes	$-2.1 \cdot 10^{-37}$	$-2.1 \cdot 10^{-33}$	$3.8 \cdot 10^{-34}$	$8.3 \cdot 10^{-36}$
<u>u quark EDM</u>						
Gluino loop : -						
α_s	1	yes	0	0	$4.5 \cdot 10^{-26}$	$-1.8 \cdot 10^{-30}$
Neutralino-like loop : -						
g^2	1	yes	~ 0	0	$2.6 \cdot 10^{-27}$	$-1.4 \cdot 10^{-31}$
$g \cdot y_u$	2	no	~ 0	0	$2.1 \cdot 10^{-28}$	$5.3 \cdot 10^{-31}$
y_u^2	4	yes	~ 0	0	$1.3 \cdot 10^{-37}$	$4.0 \cdot 10^{-41}$
Chargino-like loop : -						
$g \cdot y_u$	5	no	~ 0	~ 0	$-1.3 \cdot 10^{-27}$	$-3.2 \cdot 10^{-30}$
$y_u \cdot y_d$	7	RPV	$-7.6 \cdot 10^{-36}$	$-7.6 \cdot 10^{-33}$	$3.2 \cdot 10^{-34}$	$6.4 \cdot 10^{-34}$
$y_u \cdot \chi'_{ijk}$	7	RPV	$9.7 \cdot 10^{-36}$	$9.7 \cdot 10^{-33}$	$-5.1 \cdot 10^{-34}$	$-6.4 \cdot 10^{-34}$
<u>Neutron EDM</u>						
from Gluino loop :			$1.8 \cdot 10^{-27}$	$1.8 \cdot 10^{-24}$	$-1.0 \cdot 10^{-25}$	$-1.4 \cdot 10^{-28}$
from Chargino-like loop :			$4.3 \cdot 10^{-27}$	$4.3 \cdot 10^{-24}$	$2.5 \cdot 10^{-26}$	$2.4 \cdot 10^{-28}$
from Neutralino-like loop :			$-2.9 \cdot 10^{-28}$	$-2.9 \cdot 10^{-25}$	$-8.6 \cdot 10^{-28}$	$-2.0 \cdot 10^{-29}$
<u>TOTAL :</u>			$5.8 \cdot 10^{-27}$	$5.8 \cdot 10^{-24}$	$-7.8 \cdot 10^{-26}$	$8.0 \cdot 10^{-29}$

FIGURES

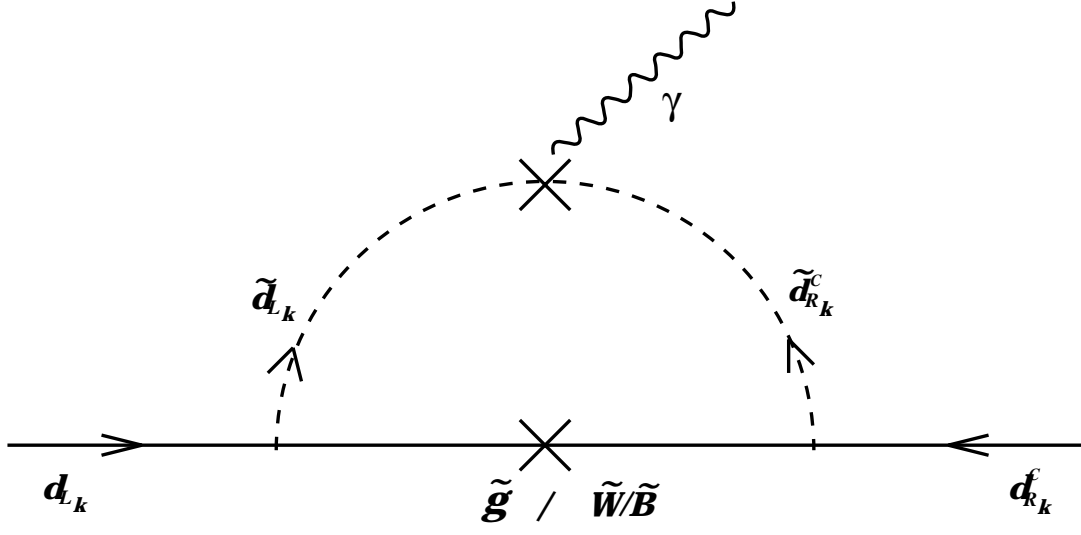


FIG. 1. Diagram for d quark EDM from gaugino loop.

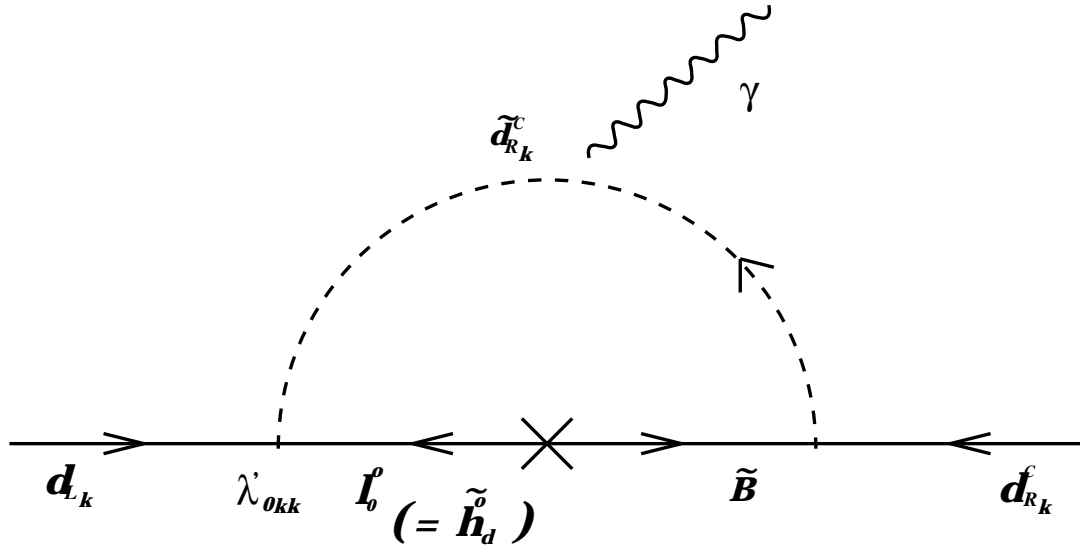
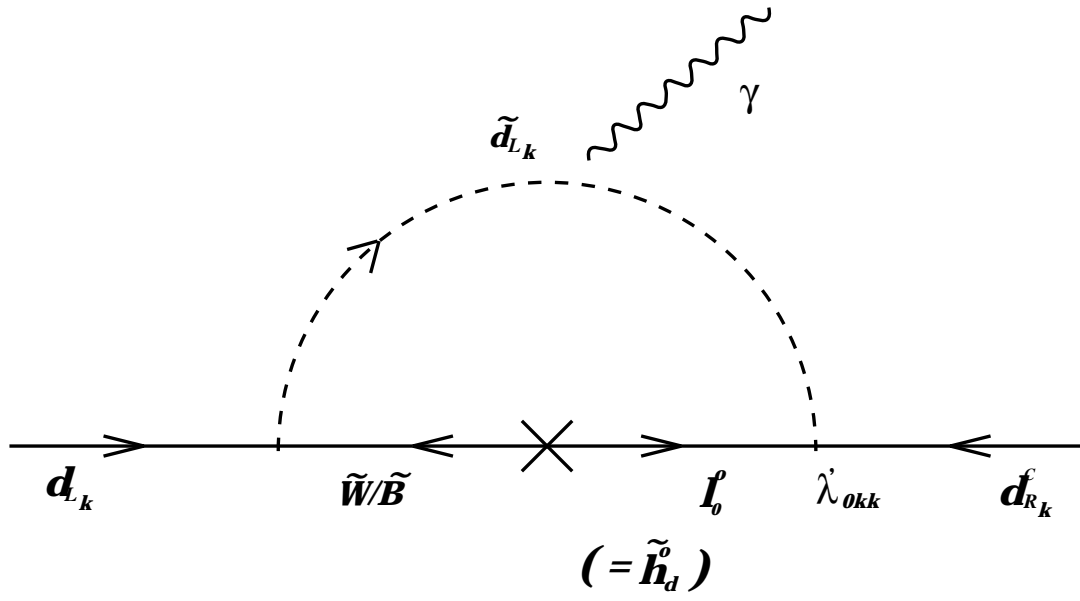


FIG. 2. Diagrams with neutral gaugino/higgsino mixing for d quark EDM.

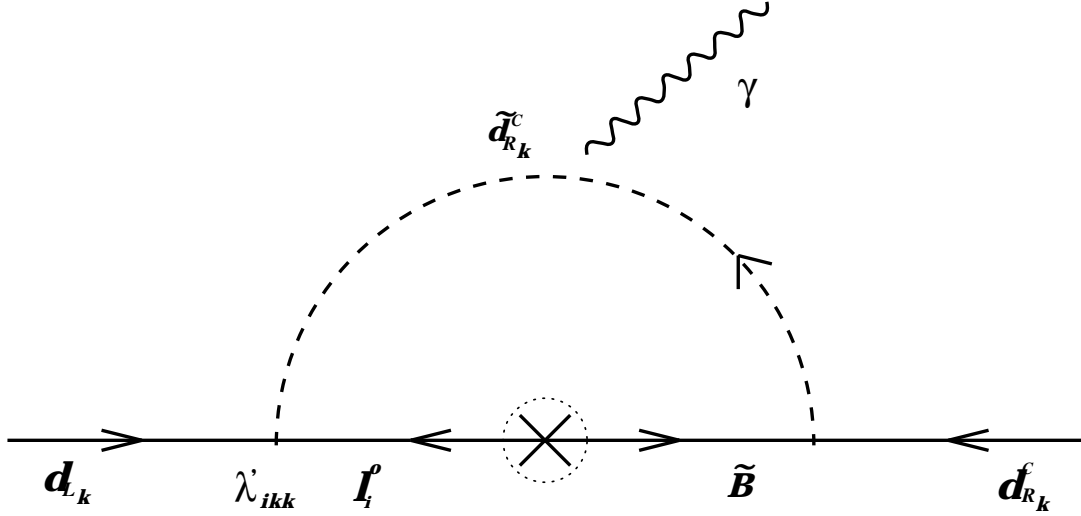
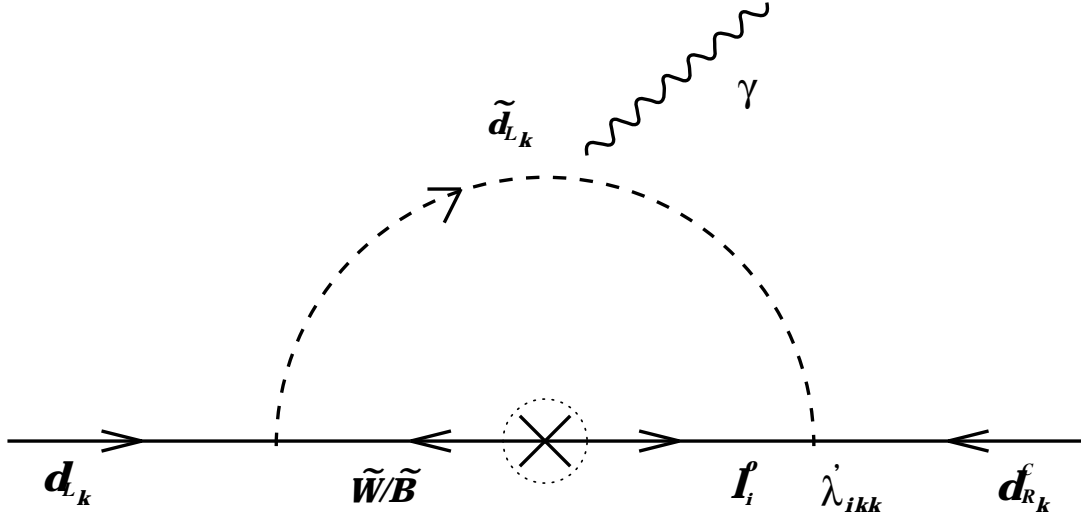


FIG. 3. R-parity violating neutralino-like loop diagrams for d quark EDM. Naive electroweak state analysis suggests that such a diagram is proportional to a vanishing l_k^0 -gaugino mass mixing.

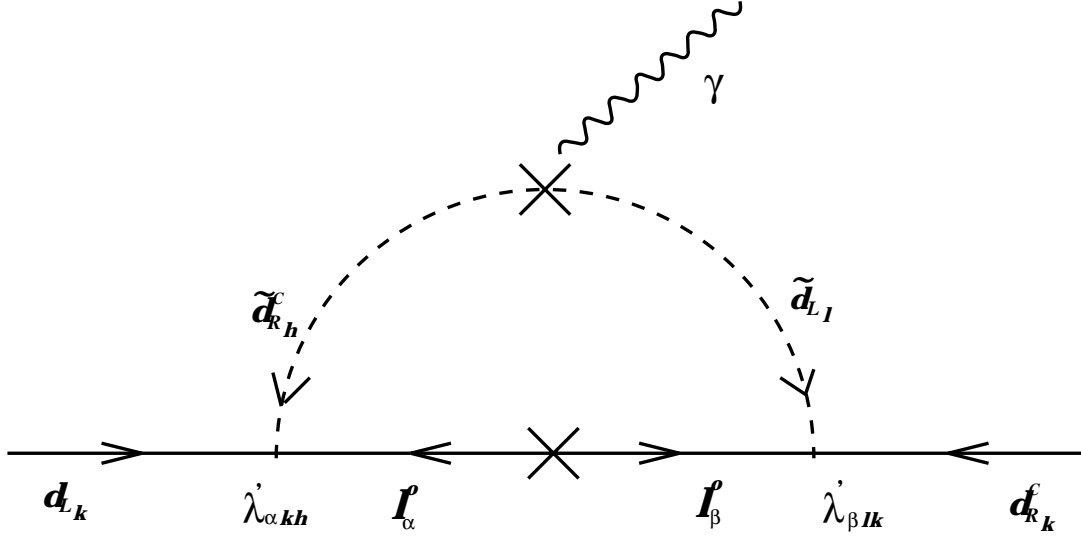


FIG. 4. Diagram for d quark EDM suggesting involvement of Majorana masses among the l_α^0 or “neutrinos”.

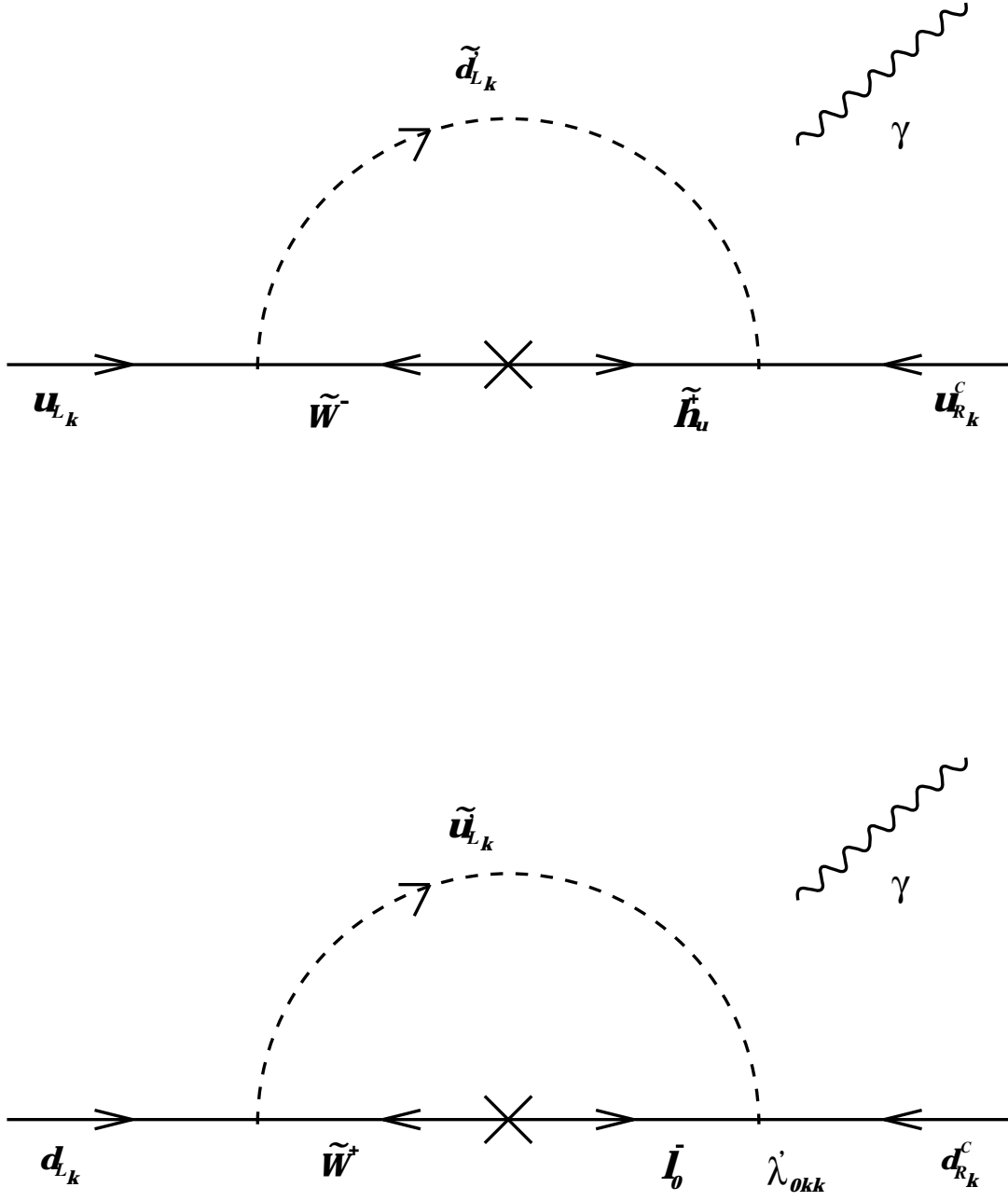


FIG. 5. Diagrams for u and d quark EDM's with charged gaugino/higgsino mixing.

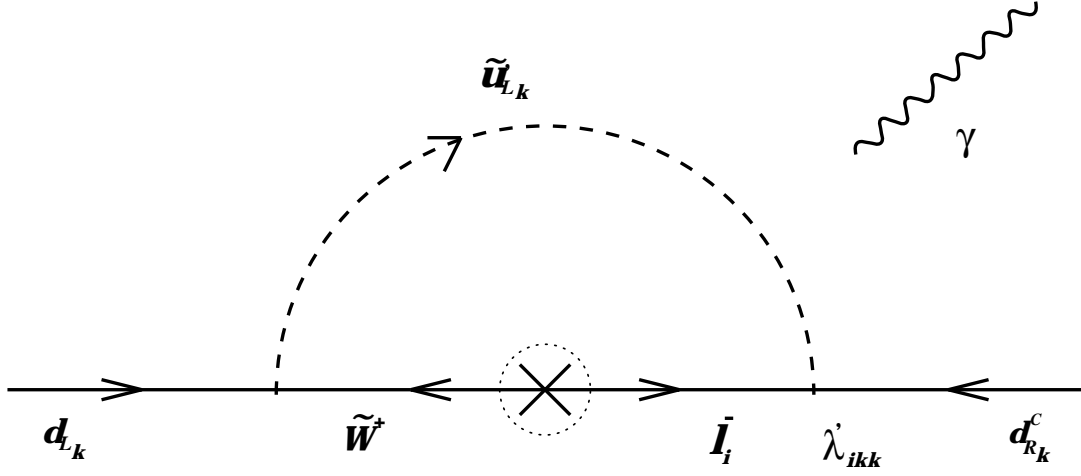


FIG. 6. R-parity violating chargino-like loop diagram for d quark EDM. Naive electroweak state analysis suggests that the diagrams are proportional to the vanishing $l_k^- \tilde{W}^+$ mass term.

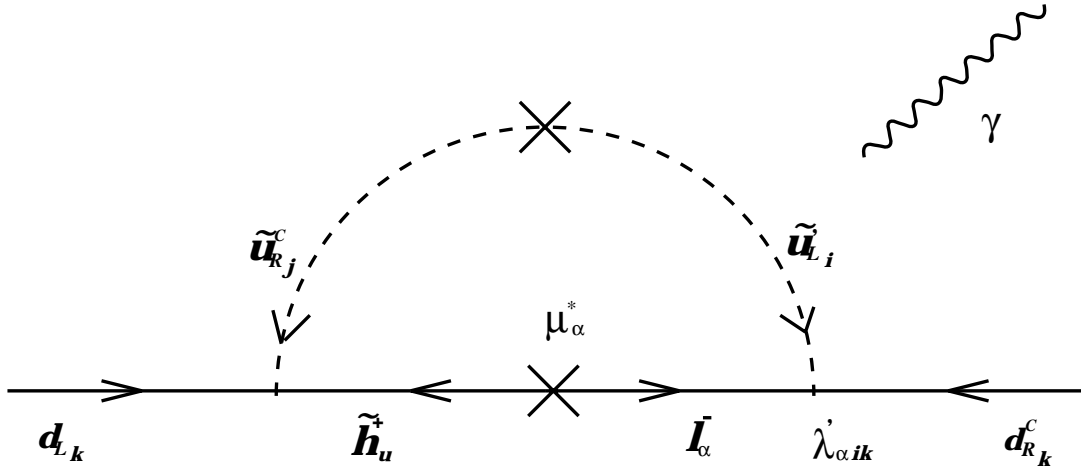
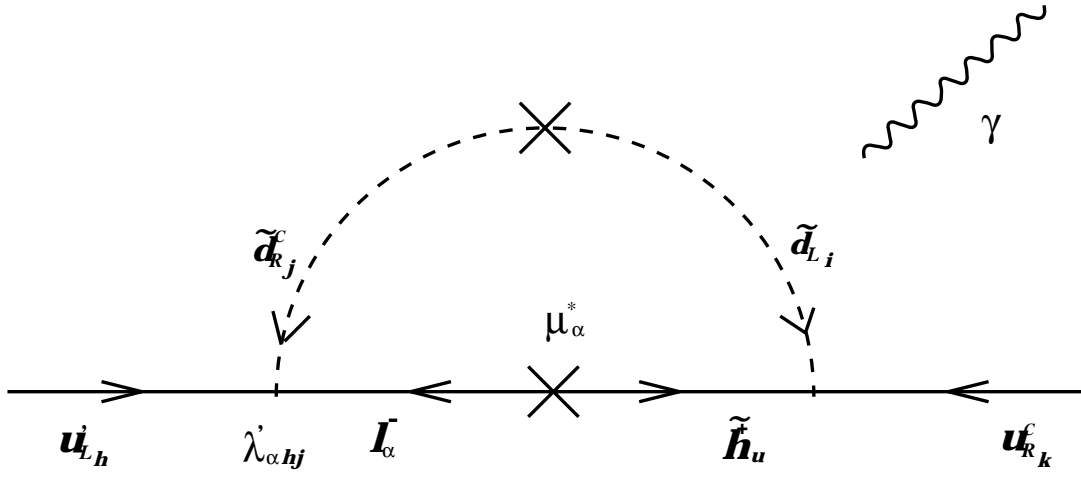


FIG. 7. Diagrams for u and d quark EDM's with a μ_α mass insertion.

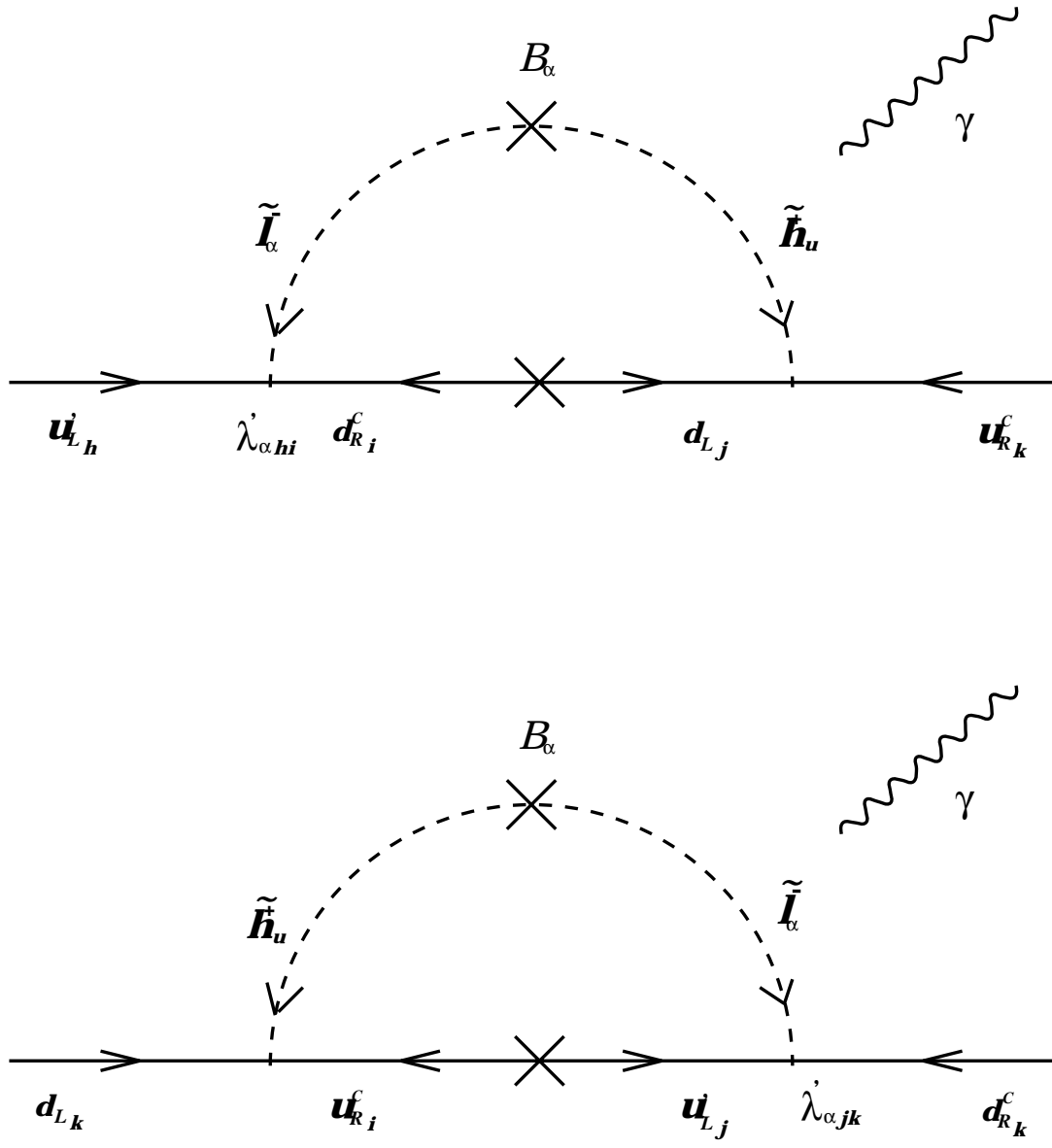


FIG. 8. Diagrams for u and d quark EDM's with a B_α scalar mass insertion.

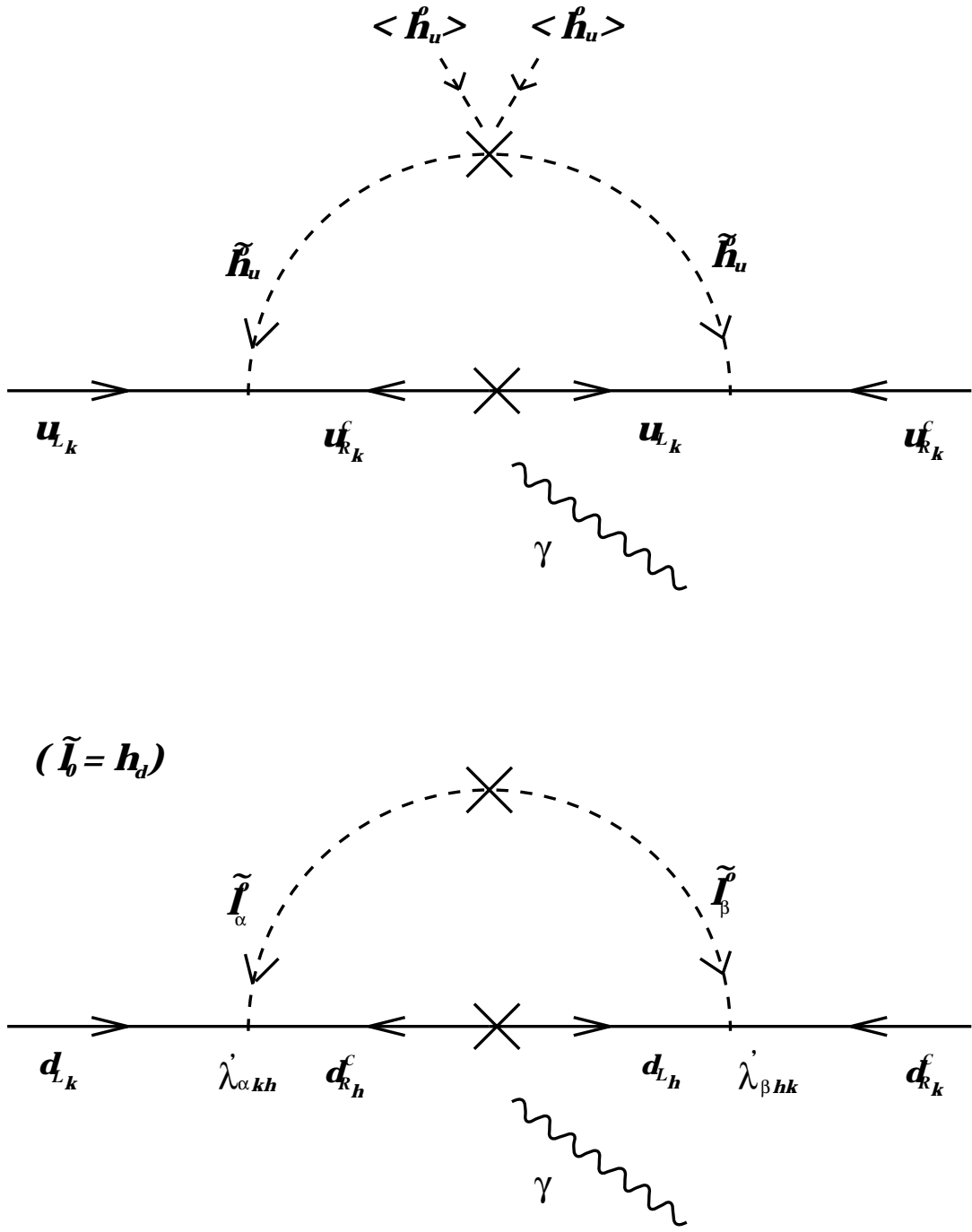


FIG. 9. Diagrams with a Majorana-like scalar mass insertion for u and d quark EDM's. For the u quark case, the direct Majorana-like h_u mass insertion is explicitly shown. For the d quark case, the corresponding direct h_d mass insertion is obvious, for $\alpha = \beta = 0$; for α and/or β nonzero, the naive direct result from the diagram would vanish, due to the vanishing VEV's.

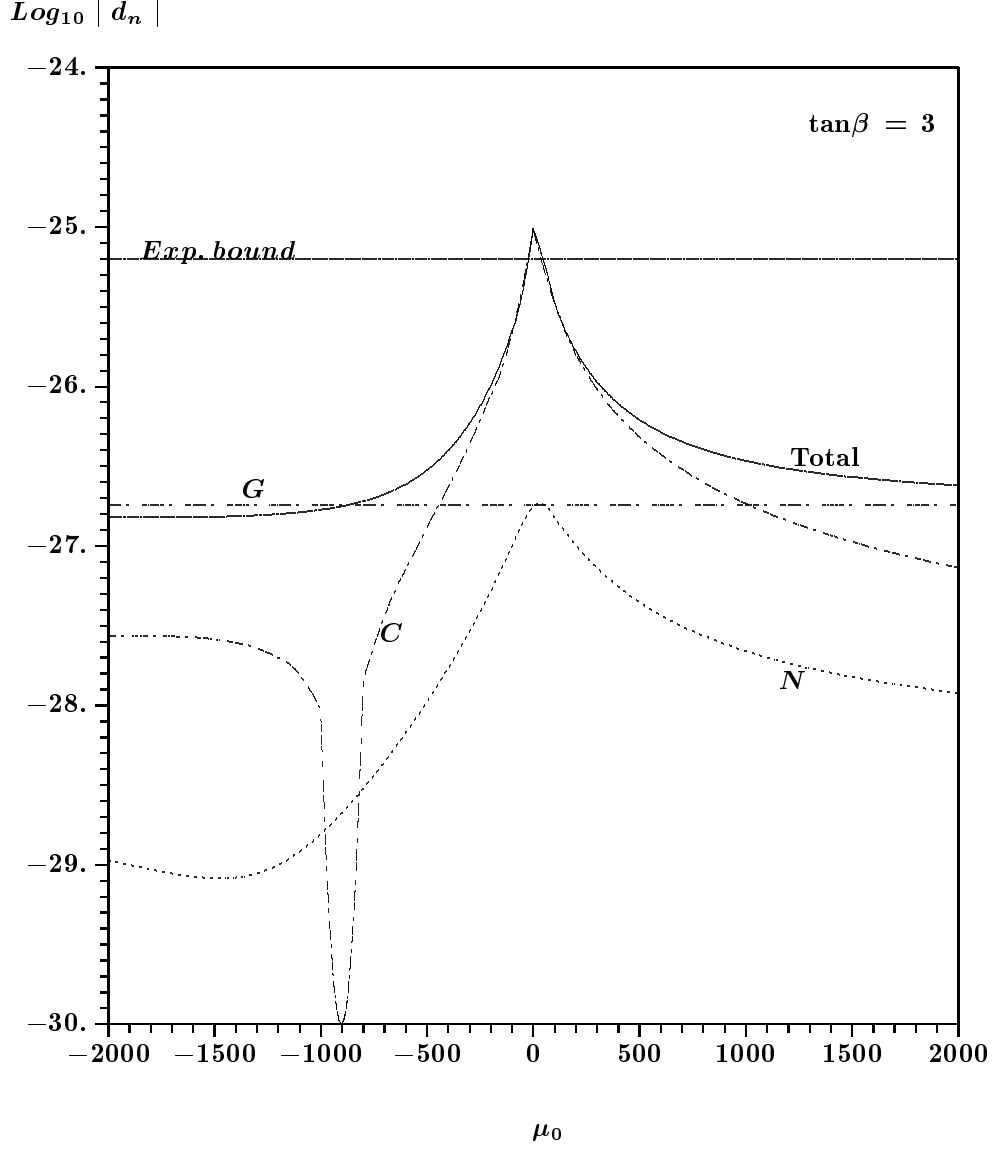


FIG. 10. Logarithmic plot of (the magnitude of) the RPV neutron EDM result for μ_0 value between ± 2000 GeV, with the other parameters set at the same values as Case A in Table 1. The lines marked by G , C , N , and “Total” gives the complete gluino, chargino-like, neutralino-like, and total (*i.e.* sum of the three) contributions, respectively. Note that the values of the N contributions and those of the C line for $\mu_0 < -900$ GeV are negative.

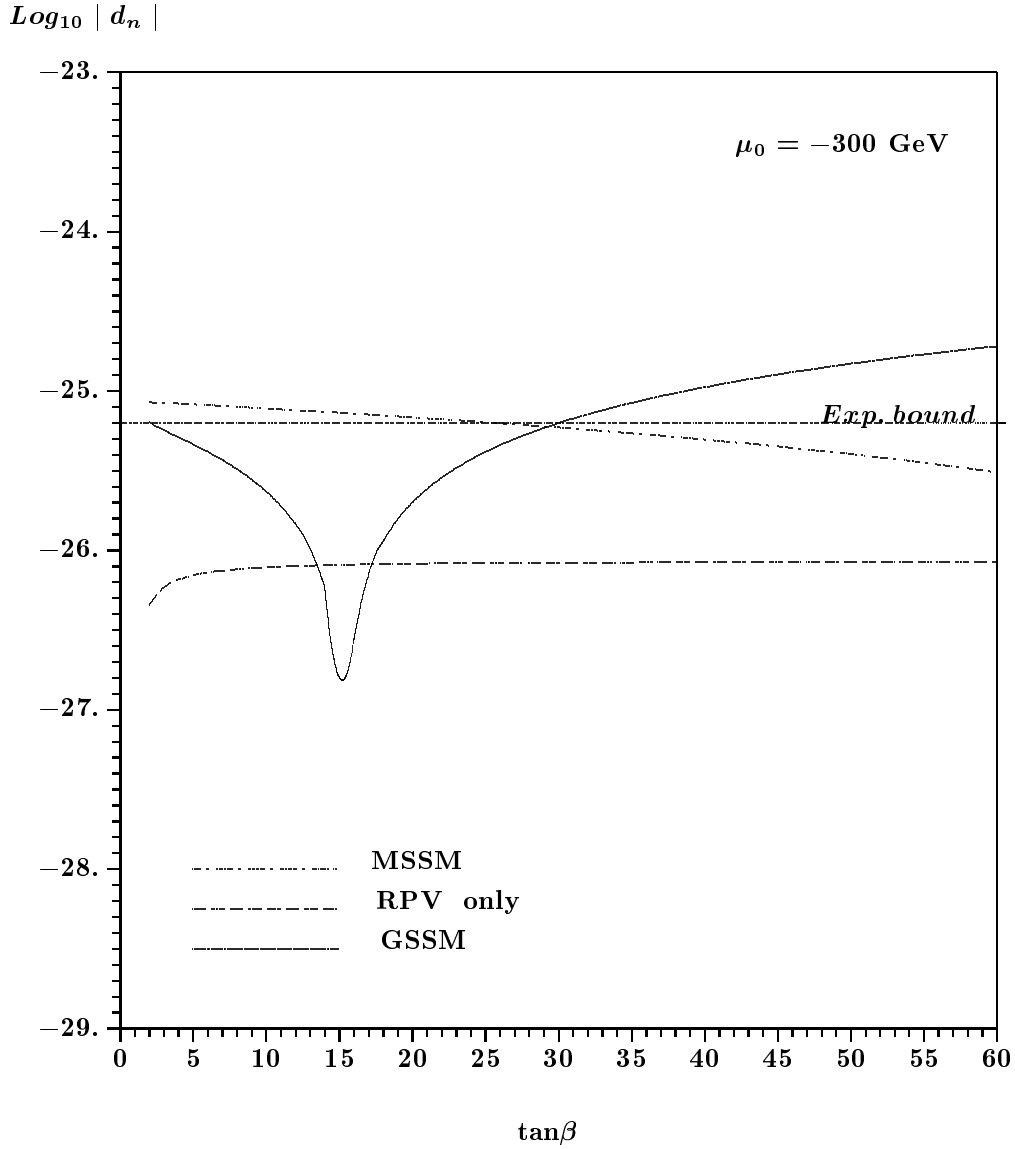


FIG. 11. Logarithmic plot of (the magnitude of) the neutron EDM result verses $\tan\beta$. We show here the MSSM result, our general result with RPV phase only, and the generic result with complex phases of both kinds. In particular, the A and μ_0 phases are chosen as 7° and 0.1° respectively, for the MSSM line. They are zero for the RPV-only line, with which we have a phase of $\frac{\pi}{4}$ for λ_{311} . All the given nonzero values are used for the three phases for the generic result (from our complete formulae) marked by GSSM. Again, the other unspecified input parameters are the same as for Case A of Table 1.

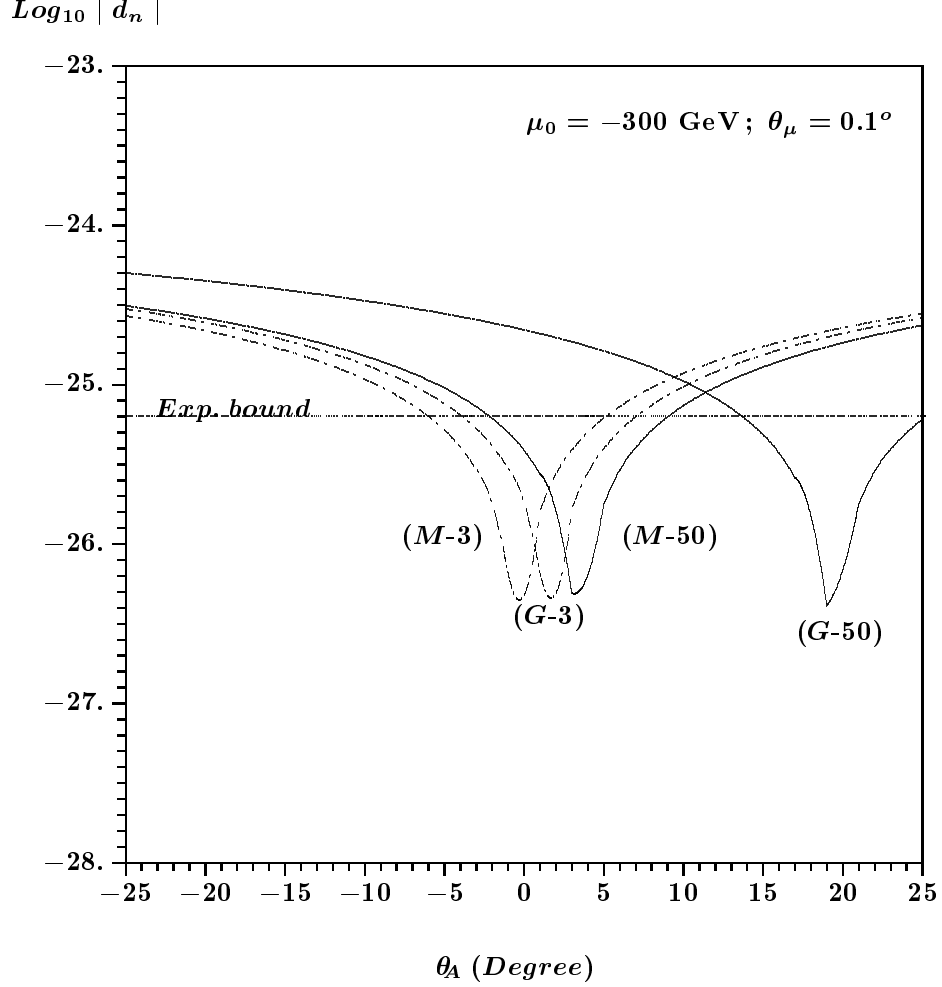


FIG. 12. Logarithmic plot of (the magnitude of) the neutron EDM result versus θ_A (the complex phase for the A parameter). The four line shown are characterized by the $\tan\beta$ values (3 or 50) used and whether it is for our GSSM result (again with a phase of $\frac{\pi}{4}$ for χ'_{3H}) — marked by G , or the result for MSSM — marked by M . Again, the χ'_{3H} phase is set at $\frac{\pi}{4}$ for the G lines, and the μ_0 phase at 0.1° for all; the other unspecified input parameters are the same as for Case A of Table 1.

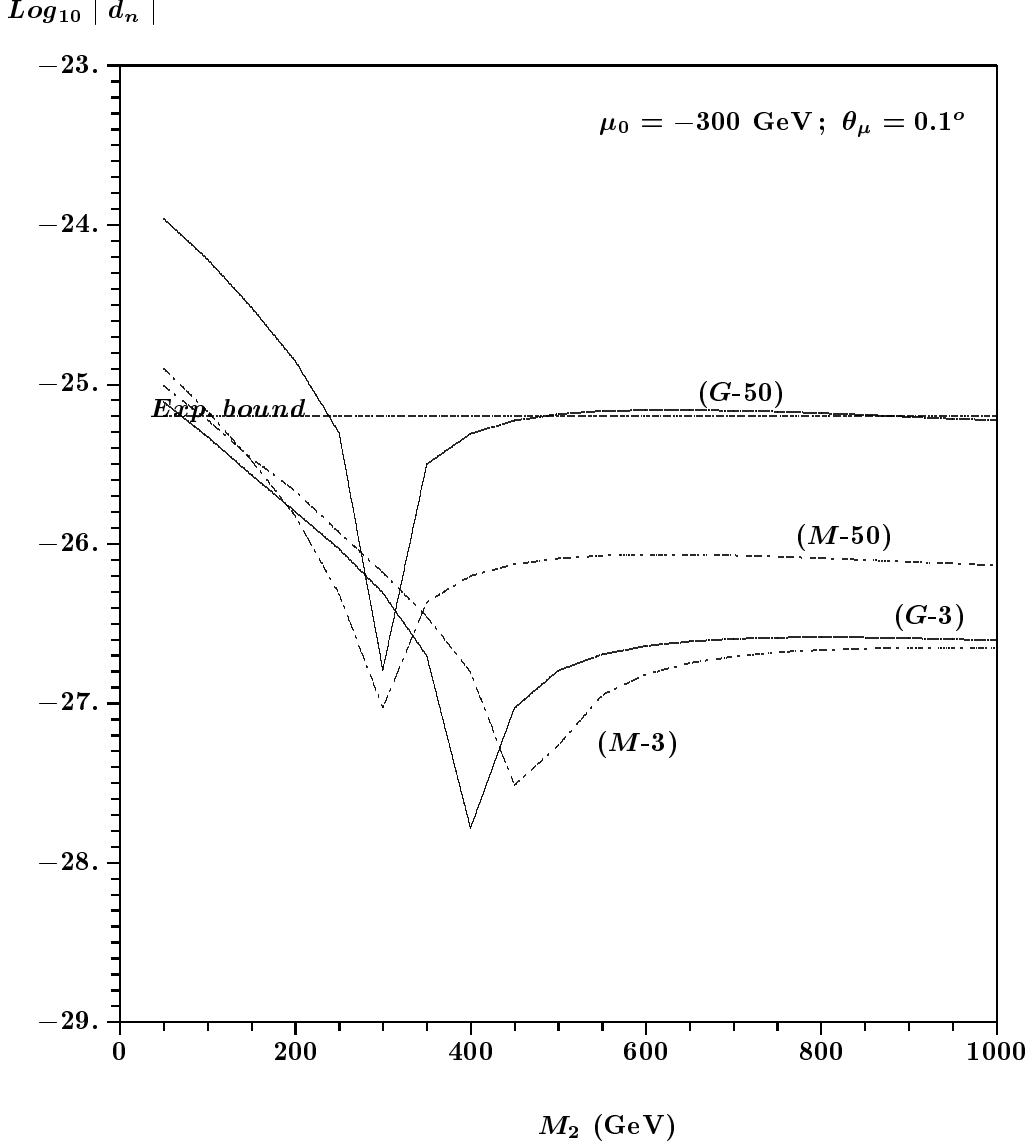


FIG. 13. Logarithmic plot of (the magnitude of) the neutron EDM result versus M_2 . The four lines correspond to the four cases of Fig. 12 each with θ_A set at the dip location, *i.e.* $G-3$ for GSSM at $\tan\beta = 3$ with $\theta_A = 2^\circ$, $M-3$ for MSSM at $\tan\beta = 3$ with $\theta_A = -1^\circ$, $G-50$ for GSSM at $\tan\beta = 50$ with $\theta_A = 20^\circ$, $M-50$ for MSSM at $\tan\beta = 50$ with $\theta_A = 3^\circ$. All other unspecified input parameters are the same as for Fig. 12.